THÈSE

présentée

A L'UNIVERSITÉ DE PARIS VII

pour obtenir

LE GRADE DE DOCTEUR ES_SCIENCES PHYSIQUES

par

Georges GIRARDI

Première Thèse

CONTRIBUTION A L'ETUDE PHENOMENOLOGIQUE

DES REACTIONS MESON-BARYON A HAUTE ENERGIE

Soutenue le 4 Février 1976 devant la Commission d'Examen

Président Mr. B. DIU

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INTRODUCTION GENERALE

La matière de cette thèse est composée d'articles déjà publiés dans des revues, et donc rédigés en anglais ainsi que d'une partie [chapitre III] "inédite", en français. Les travaux sont présentés dans l'ordre chronologique - mis à part le chapitre III - ce qui permet de suivre l'évolution des conceptions de l'auteur.

Nous avons séparé l'exposé en deux parties, correspondant aux deux approches utilisées dans l'étude des réactions à 2 corps. Chacune de ces deux parties est précédée d'une introduction ou plutôt d'un sommaire des résultats présentés. Une telle présentation qui peut conduire à des redites, a par contre l'avantage de faciliter l'accès aux détails des articles pour le lecteur motivé.

La première partie est consacrée à l'élaboration d'un nouveau modèle phénoménologique. Dans cette construction, nous faisons appel à un large éventail de concepts théoriques, tels la dualité, les diagrammes de quarks, la symétrie SU(3) etc ... Le modèle ainsi élaboré est confronté aux résultats expérimentaux et paraît résoudre certaines difficultés que d'autres modèles ne peuvent éviter.

Dans la seconde partie l'approche adoptée est plus typiquement phénoménologique. En effet, elle est consacrée à des analyses en amplitudes, c'est-à-dire que l'on interroge l'expérience, en évitant - autant que possible - de faire appel à des concepts théoriques par trop particuliers. Les divers résultats obtenus permettent de percevoir certaines régularités dans les amplitudes, qui sont la manifestation de lois physiques sous-jacentes encore mal comprises. Cette étude nous a permis de proposer une systématique pour les amplitudes de diffusion à deux corps. Cette systématique qui simplifie considérablement la formulation des modèles phénoménologiques a déjà été utiliséepour étudier certains processus avec succès. D'autres études sont actuellement en cours, dont les résultats semblent prometteurs. Cependant le problème théorique reste posé, au delà de la constatation de cette systématique, quelle en est l'origine et dans quelle mesure est-elle exacte ? Nous pensons que l'intérêt de ce travail

est de présenter des éléments qui pourront permettre de donner une réponse et également de réunir les deux approches utilisées en phénoménologie des interactions fortes : la construction explicite de modèles et l'analyse en amplitudes. En particulier, nous montrons que ces deux points de vue conduisent à des résultats concordants, ce qui est un indice de la pertinence des images physiques utilisées et permet d'aller plus avant dans la compréhension des interactions hadroniques à haute énergie.

1ère PARTIE

CONSTRUCTION D'UN MODELE PHENOMENOLOGIQUE

A COUPURES DE REGGE

:-:-:-:

INTRODUCTION

Le premier chapitre traite des problèmes théoriques des coupures de Regge. Il constitue la matière d'un cours publié fait à l'Ecole d'Eté de Physique des Particules de Gif-sur-Yvette en 1972. La présence de coupures dans le plan du moment angulaire complexe y est justifiée (problème des pôles fixes et de l'unitarité). Nous abordons ensuite le calcul explicite des coupures engendrées par l'échange de deux Reggeons dans le langage des vertex de Gribov. L'accent est mis sur le caractère non plan des diagrammes qui contribuent effectivement aux termes de coupures ainsi que sur l'analogie avec les termes de double diffusion de la théorie de Glauber.

Ces remarques conduisent à des règles précises pour la construction de diagrammes de quarks représentant l'échange de coupures de Regge. Tous ces résultats sont exploités dans le chapitre II où nous proposons un modèle à coupures de Regge utilisant les diagrammes de quarks et leurs règles de sélection. Ce travail a été motivé par les difficultés rencontrées par les modèles à pôles de Regge avec ou sans la correction de la coupure d'absorption. Le plus connu de ces problèmes est celui de la polarisation dans la réaction avec échange de charge $\pi^- p \to \pi^o n$.

Bien qu'on ne puisse dire par conséquent que les pôles de Regge constituent une théorie des interactions hadroniques, ils projettent néanmoins quelque lumière sur la physique des interactions fortes. La plupart des idées physiques dans ce domaine ont d'ailleurs émergé d'analyses phénoménologiques des données expérimentales, formulées en termes d'échange de pôle de Regge. Cela a conduit aux concepts de dualité et de dégénérescence d'échange et l'introduction des corrections d'absorption a permis d'affiner l'étude des liens entre pôles de Regge et unitarité. La dégénérescence d'échange donne des relations entre les trajectoires et entre les couplages de pôles de Regge différents, relations assez bien vérifiées. Cependant cette "symétrie" (comme beaucoup d'autres !) n'est pas exacte et doit être brisée. Cette brisure est particulièrement évidente dans les réactions avec échange de charge ($K^- p \to \bar{K}^o n$ et $K^+ n \to K^o p$) ou d'hypercharge ($\pi N \to K\Sigma/\Lambda$ et $\bar{K}N \to \pi\Sigma/\Lambda$). La dégénerescence d'échange prévoit en effet que les réactions reliées par croisement s-u

(renversement de lignes) telles que $\pi^+p \to K^+\Sigma^+$ et $K^-p \to \pi^-\Sigma^+$ doivent avoir même distribution angulaire. Cette prédiction est contredite par l'expérience, les distributions angulaires diffèrent fortement, sinon en structure, du moins en grandeur. Cette violation de la symétrie par renversement de lignes est encore plus frappante dans les analyses en amplitudes. Elles indiquent en effet que les amplitudes non flip de signature négative (échange de vecteur : $\rho,\omega,K^*(892)$) sont périphériques, c'est-à-dire, ont une structure en transfert décrite par une fonction $J_o(R\sqrt{-t})$ (pour la partie imaginaire). La fonction de profil dans l'espace du paramètre d'impact est de la forme $\delta(b^2-R^2)$. Tout se passe comme si la cible était un anneau de rayon $R \sim 1$ fm, l'essentiel de la diffusion se faisant à la périphérie du hadron. Une telle propriété est bien expliquée par les modèles d'absorption. Par contre la structure des amplitudes de signature positive (échange de tenseur $A_2, f_o, K^{**}(1420)$) est plutôt du type central, c'est-à-dire correspond à une fonction de profil de la forme $\theta(R^2-b^2)$: le hadron se comporte comme un disque plus ou moins uniforme. On est donc confronté au problème suivant : comment briser la dégénérescence d'échange dans les amplitudes sans perdre les principaux résultats de la dualité et surtout la simplicité et l'économie de paramètres impliquées par ses conséquences sur la dégénérescence des trajectoires ?

La solution que nous proposons est que la dégénérescence d'échange, exacte au niveau des termes de pôle, se trouve brisée par les contributions des coupures de Regge. Cette idée est confirmée par le fait que la brisure de la symétrie par renversement de lignes disparaît rapidement quand l'énergie augmente. L'effet doit donc être attribué à des singularités de Regge de trajectoires plus basses que celles des pôles. Notre point de vue est que les trajectoires secondaires associées à ces singularités sont celles de coupures de Regge. D'autres auteurs introduisent des trajectoires filles pour rendre compte des anomalies observées. Le problème dans ce cas là est l'absence d'association pôle de Regge-particule : on ne connait pas de trajectoires fille portant des particules physiques - sauf peut être le $\rho'(1600)$ mais il est surtout couplé à 4π - . De ce fait, les modèles à trajectoires filles nous paraissent être essentiellement une paramétrisation utile des données expérimentales.

Pour notre part, nous considérons que l'échange d'un pôle de Regge est le terme de Born d'une théorie complète et que, en tant que tel, il rend compte des faits essentiels. Ainsi la présence d'un pic à l'avant (resp. à l'arrière) avec une section efficace substantielle dans une distribution angulaire est strictement corrélée à la non exoticité des nombres quantiques de la voie t (resp.u). Vouloir aller plus loin dans la description des phénomènes nécessite de prendre en compte les termes d'échanges multiples, les itérations du terme de Born en particulier. On voit dans le chapitre I qu'à ce niveau on introduit a priori beaucoup d'inconnues, telles les vertex de Gribov qui sont des intégrales sur la région physique des réactions Reggeon-particule → Reggeon-particule. Afin de ne pas donner trop de flexibilité au modèle par un nombre excessif de paramètres, nous allons imposer des règles de calcul précises. Notre point de départ, notre germe, est un pôle de Regge, vérifiant la dégénérescence d'échange, construit par un diagramme de quarks. Nous engendrons les termes successifs d'échanges multiples à partir de ce germe. Nous retrouvons de cette façon l'absorption habituelle et de nouveaux termes dont l'effet est de corriger les insuffisances du modèle d'absorption usuel. En particulier ces termes fournissent un mécanisme de brisure de la dégénérescence d'échange, un des termes essentiels étant la coupure Reggeon Pomeron Reggeon (l'ordre est important ! RPR). L'accord obtenu avec les résultats expérimentaux sur la diffusion $\pi N \to \pi N$ et $KN \to KN$ est très bon avec un nombre réduit de paramètres. Ces paramètres sont ajustés en utilisant les données expérimentales sur $\pi N \to \pi N$. A l'aide de la symétrie SU(3) nous prédisons $\pi N \to \pi N$ et $KN \to KN$, en particulier nous obtenons :

$$\frac{d\sigma}{dt}(K^+ n \to K^o p) > \frac{d\sigma}{dt}(K^- p \to \bar{K}^o n) .$$

L'effet de brisure de dégénérescence d'échange va dans le bon sens mais n'a pas la bonne grandeur. Une discussion complète du modèle a été donnée par R. Peschanski dans sa thèse de Doctorat d'Etat. Encouragés par ces résultats, nous avons appliqué le modèle RPR pour décrire les réactions avec échange d'hypercharge ($\pi N \to K\Sigma/\Lambda$ et $\bar{K}N \to \pi\Sigma/\Lambda$) dans lesquelles la brisure de la dégénérescence d'échange est particulièrement forte. Le traitement de ces réactions est présenté dans le chapitre III.

Nous utilisons la symétrie SU(3) afin de calculer tous les couplages à partir des valeurs trouvées dans π-N. Pour ce faire nous sommes conduits à introduire des hypothèses simplificatrices, peut être injustifiées. Les résultats présentées sont donc en fait des prédictions, l'accord semi-quantitatif avec l'expérience doit être considéré comme satisfaisant étant données les contraintes imposées au modèle. On obtient une description correcte pour la brisure de la dégénérescence d'échange, qui est un problème pour beaucoup de modèles ayant des paramètres ajustables. De la même façon, le modèle prédit un ordre de grandeur raisonnable pour les réactions avec voie t exotique telles $\pi^- p \to K^+ \Sigma^-$ et $K^- p \to \pi^+ \Sigma^-$. Dans ces réactions, il y a un double échange de charge dans la voie t et seules les coupures à au moins deux Reggeons échangés contribuent. On a ici un test du calcul des coupures Regge-Regge et une mesure des vertex de Gribov. Les amplitudes de diffusion obtenues pour les échanges vecteur et tenseur vont dans le sens de la systématique proposée dans la 2ème partie.

CHAPITRE I

I . ASPECTS THEORIQUES DES COUPURES DE REGGE

I .1. Coupures de Regge et pôles fixes

I .1.1. Introduction

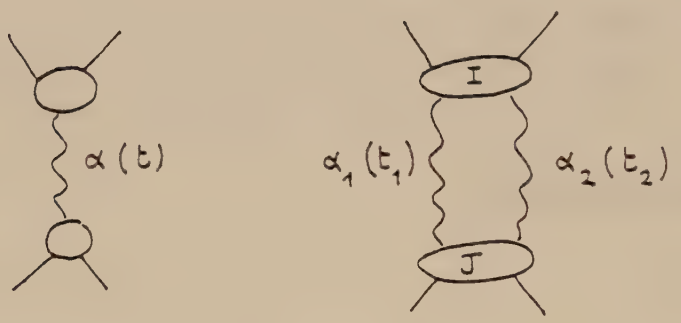
Nous avons vu dans les cours précédents que les réactions à 2 corps et à quasi-2 corps, pouvaient être décrites à haute énergie par l'échange de pôles dans le plan complexe du moment angulaire [pôles de Regge], échange qui a lieu dans la voie croisée de la réaction considérée.

Nous avons constaté par l'analyse des résultats expérimentaux et de leur description phénoménologique qu'un modèle à pôles de Regge n'était pas entièrement satisfaisant. On peut résoudre certains problèmes en introduisant de nouvelles trajectoires qui ne portent aucune particule physique, telles le ρ' et le P'''.

Il semble plus naturel de considérer que le terme de pôles n'est que le premier ordre d'une théorie complète et donc de chercher les termes d'ordre supérieur. Si l'on fait le parallèle avec la théorie

du potentiel, le terme de pôles est le terme de Born, les autres termes correspondent aux itérations du potentiel et donc à des échanges multiples (fig. I ,1 a,b).

Ainsi le terme du second ordre (fig. I ,I,b) comporte l'échange de 2 "Reggeons". Les boules I et J peuvent être des processus compliqués, car nous n'avons pas, a priori, de prescription pour itérer les échanges de Reggeons.



(a) (b)

"Terme de Born" Terme du 2nd ordre

Fig. I .1

De façon générale, ce sont des termes d'échanges multiples qui donnent lieu à des coupures dans le plan du moment angulaire.

Nous verrons que I et J doivent avoir des propriétés particulières pour que les termes de double diffusion contribuent à l'amplitude physique.

I .1.2. Les pôles fixes

D'un point de vue purement théorique, il n'y a aucun obstacle à la présence de coupures dans le plan l . Mieux, les coupures permettent de résoudre le problème des pôles fixes non sens.
Ceux-ci apparaissent quand on prolonge la formule de Froissart-Gribov à $l < 0$ (valeurs non sens) pour $l = -1, -2, \ldots, -n$, à cause des $Q_l(\cos\Theta)$.

Or si on écrit l'unitarité élastique dans la voie t (I .2)

$$a_l(t) = \frac{1}{\pi}\int_{z_0}^{\infty} A_s(s,t)\ Q_l(z)\,dz + \frac{1}{M}\int_{z_0}^{\infty} A_\mu(u,t)\ Q_l(z')\,dz' \qquad (I.1)$$

$$a_l^*(t)^* - a_l(t) = \sqrt{t - 4m^2}\ |a_l(t)|^2 \qquad (I.2)$$

On a un pôle simple à gauche et un pôle double à droite. Un pôle fixe est incompatible avec l'unitarité, mais pas un pôle de Regge puisqu'il se déplace avec t et qu'au seuil la trajectoire a une partie imaginaire.

Une méthode astucieuse pour éliminer ces singularités génantes est de les faire passer dans un feuillet non physique du plan. Pour cela, il suffit qu'il y ait dans le plan l un point de branchement mobile passant par le pôle à $t = t_0$ valeur au seuil, le pôle est alors envoyé dans un autre feuillet, on dit que la coupure recouvre [en anglais shields] le pôle. Le mécanisme est plus clair sur un exemple.

Considérons l'amplitude d'onde partielle suivante :

$$a(l,t) = - \frac{(t - m^2)^l \; \varphi(l,t)}{\dfrac{\varphi(-1,t)}{\sqrt{\dfrac{l - \alpha_c'(t)}{\alpha_c'}}} + \dfrac{(t - 4m^2)^{l+1} \; \varphi(l,t)}{\sqrt{\dfrac{-1 - \alpha_c(t)}{\alpha_c'}}}} \qquad (I.3)$$

où $\varphi(l,t)$ est analytique en l et t et
$\alpha_c(t) = -1 + \alpha_c'(t - 4m^2)$

On peut vérifier que $a(l,t)$ a bien un pôle fixe à $l = -1$, une coupure mobile passant par $l = -1$ à $t = 4m^2$ et qu'elle vérifie la condition d'unitarité avec le bon comportement au seuil.

Ce mécanisme a été utilisé par Mandelstam [1] pour justifier l'existence de coupures de Regge.

Il y a plusieurs méthodes pour étudier les singularités que donnent des "graphes" comme celui de la fig. 1,b. Ici nous utiliserons les techniques des graphes de Feynman. Etant donné la valeur importante de leur constante de couplage, une théorie perturbative des champs pour les interactions fortes n'a aucune chance d'être satisfaisante. Toutefois, il est admis que par l'étude des intégrales de Feynman, on peut atteindre les propriétés d'analyticité des amplitudes de diffusion. [2]

C'est la philosophie de la théorie de la matrice S où les équations de Landau sont une systématisation de l'étude des singularités des graphes dans une théorie en ϕ^3 - tous les vertex comportent 3 particules. Il en est de même pour l'analyticité de Mandelstam c'est une hypothèse non démontrée, les régions où l'analyticité est

est effectivement démontrée, sont restreintes. [3]

C'est dans cet esprit que nous étudierons le terme à 2 Reggeons (fig. 1,b) de façon à définir les vertex de Gribov et à déduire les règles de sélection qu'imposent I et J. [4]

Nous étudierons ensuite le cas où I et J ne comportent qu'une particule intermédiaire. Cette simplification nous ramènera au calcul des modèles d'absorption et à la mise en évidence de coupures de Regge. Les calculs détaillés sont donnés dans l'appendice B. [5]

Dans la dernière partie, nous faisons le parallèle : coupures de Regge - Théorie du Deuteron de Glauber, ce qui nous permettra d'introduire la notion de nature composite des hadrons (quarks, partons).

I .2 Les vertex de Gribov [4]

I .2.1. Comportement de Regge et graphes de Feynman

Nous rappelons simplement qu'on peut obtenir un comportement de Regge en sommant certains types de graphes : par exemple la série infinie d'échelles ci-dessous donne un comportement asymptotique en $S^{\alpha(t)}$

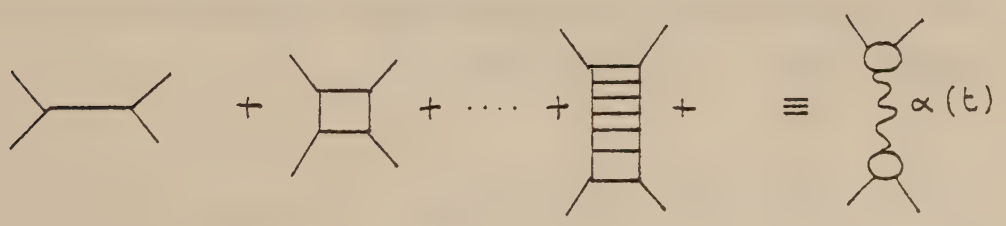


Fig. I .2

Ainsi un terme du second ordre comme celui de la figure 1,b est obtenu par sommation de la série suivante :

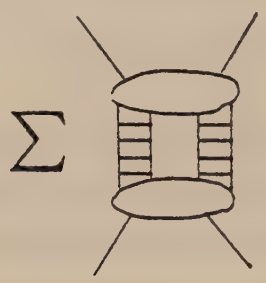


Fig. I .3

Dans le calcul des graphes nous exprimerons l'échange du pôle de Regge par la forme habituelle sans tenir compte de la signature.

$$R(s,t) = \gamma(t)\,\Gamma(1-\alpha(t))\,S^{\alpha(t)} \qquad (I.4)$$

Toutefois, à l'intérieur d'un graphe cette amplitude sera prise hors de sa couche de masse, il faudra supposer que l'extrapolation se fait sans gros problèmes.

I .2.2. Termes à 2 Reggeons

Considérons le graphe de la figure 4, nous étudions son comportement asymptotique pour $s = (p_1 + p_2)^2$ grand. Pour cette étude nous utiliserons le repère du cône de lumière ou repère de moment infini. Les 4 vecteurs sont repérés par 2 composantes sur des vecteurs du genre lumière l_+ et l_- avec $l_\pm = \frac{1}{\sqrt{2}}(1,0,0,\pm 1)$ et par 2 composantes dans un diplan normal au cône de lumière.

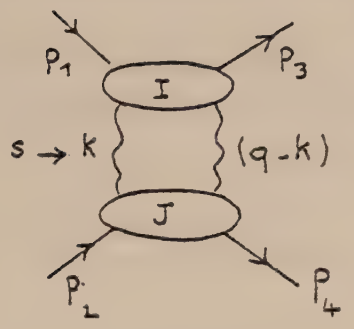


Fig. I..4

L'adoption de ce repère revient à faire le changement suivant :

$$k = (k^o, \vec{k}) \rightsquigarrow (k^+, k^-, \vec{k}_\perp) \quad \text{avec}$$

$$k^\pm = \frac{1}{\sqrt{2}} (k_o \pm k_3) = \zeta_\pm \cdot k \qquad (I .5)$$

On a donc pour le produit scalaire :

$$\left.\begin{aligned} k_i \, k_j &= (k_i^+ \, k_j^- + k_i^- \, k_j^+) - k_{i\perp} \, k_{j\perp} \\ k_i^2 &= 2 \, k_i^+ \, k_i^- - k_{i\perp}^2 \end{aligned}\right\} \qquad (I .6)$$

Si nous appelons I et J les amplitudes de diffusion Particule Reggeon $\longrightarrow$ Particule Reggeon l'expression à calculer s'écrit :

$$A(s,t) = \int I(p_1, p_3, k, q-k) \; J(p_2, p_4, k, q-k) \; R_1(k^2, s) \; R_2((q-k)^2, s) \; d^4 k$$

(I .7)

Les résidus de Regge sont inclus dans I et J, on peut donc écrire :

$$R_1 (k_2, s) = S^{\alpha_1(k^2)}$$

$$R_2 ((q - k)^2, s) = S^{\alpha_2 ((q - k)^2)}$$

(I .8)

Nous nous plaçons dans les conditions cinématiques suivantes :

S est grand, R_1 et R_2 ne contribuent que si leurs énergies respectives sont grandes ;

les moments de transfert sont de l'ordre de m^2

Ces conditions se traduisent de la façon suivante :

si k_i et k_j sont respectivement les 4 moments internes de I et de J

$$k_i^+ \sim \sqrt{S} \quad , k_i^- \sim m^2/\sqrt{S} \quad \text{et } k_{i\perp}^2 \sim m^2$$

$$k_j^+ \sim m^2/\sqrt{S} \quad , k_j^- \sim \sqrt{S} \qquad k_{j\perp}^2 \sim m^2$$

(I .9)

On peut voir cela intuitivement : on a collision de 2 particules d'impulsion p_1 et p_2 soit dans le centre de masse

$$p_1^+ \sim \sqrt{S} \text{ et } p_2^- \sim \sqrt{S} \text{ et } (p_1^- \text{ et } p_2^+) \sim m^2/\sqrt{S}$$

En définissant la rapidité par : $y = \log k^+ / k^-$ on a

$$y_i = \log S \quad \text{et} \quad y_j = - \log S .$$

On a donc 2 "paquets de particules" de rapidités grandes et opposées. Il est raisonnable de penser que chacun va garder sa direction et que les seuls échanges possibles sont de nature tranverse. Cette image se retrou-

ve pour les collisions à très haute énergie dans le modèle des partons.

Ceci a une conséquence importante car les moments de transfert des 2 Reggeons sont rattachés à des k_i et des k_j, ils n'ont pas de composantes $k^{\pm}$ et $(q-k)^{\pm}$ de l'ordre de $\sqrt{S}$ donc :

$$k^2 \sim k_\perp^2 \quad \text{et} \quad (q-k)^2 \sim (q-k)_\perp^2 \qquad (I.10)$$

Ceci montre qu'il y a peu de transfert par les Reggeons et donc que I et J sont quasiment indépendantes, on va pouvoir factoriser A(s,t). Si s_1 et s_2 sont les énergies de I et J respectivement :

$$s_1 = (p_1 + k)^2 \sim \sqrt{S}\ k^-$$
$$s_2 = (p_2 - k)^2 \sim \sqrt{S}\ k^+ \qquad (I.11)$$

soit : $$d^4k = dk^+\, dk^-\, d^2k_\perp \sim \frac{1}{S}\, ds_1\, ds_2\, d^2k_\perp$$

Il vient par (I-7), (I-10) et (I-11)

$$A(s,t) \sim \int d^2k_\perp\ R_1(k_\perp^2, s)\ R_2((q-k)_\perp^2, s)\ x$$
$$x \int dk^+\, dk^-\ I(t, s_1, (q-k)_\perp^2, k_\perp^2)\ J(t, s_2, (q-k)_\perp^2, k_\perp^2)$$
$$= \int d^2k_\perp\ S^{\alpha_1(k_\perp^2) + \alpha_2((q-k)_\perp^2) - 1} \int ds_1\ I(t, s_1, t_2, t_1)\ x$$
$$x \int ds_2\ J(t, s_2, t_2, t_1)$$
$$= \int d^2k_\perp\ S^{\alpha_1(k_\perp^2) + \alpha_2((q-k)_\perp^2) - 1}\ N_1(t, t_1, t_2)\ .\ N_2(t, t_1, t_2)$$

(I.12)

où N_1 et N_2 sont par définition les vertex de Gribov. Ce sont les intégrales des amplitudes I et J décrivant la diffusion :

$$a + R_1 \longrightarrow c + \bar{R}_2 \quad \text{et} \quad b + \bar{R}_1 \longrightarrow d + R_2 \qquad [\text{fig.I.5}]$$

L'intégration sur $d^2k_\perp$ nous donne le terme de coupure des modèles d'absorption.

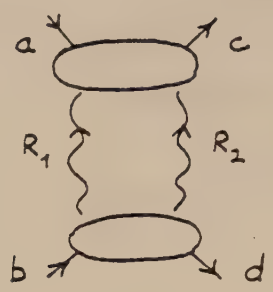


Fig. I .5

Si on admet que les amplitudes de diffusion Regge + particule $\longrightarrow$ Regge + particule ont les mêmes propriétés d'analyticité que les amplitudes particule + particule $\longrightarrow$ particule + particule, l'intégrale sur l'énergie se fait sur la région physique.

Ainsi pour $I(t, s_1, t_2, t_1)$ on a le contour C.

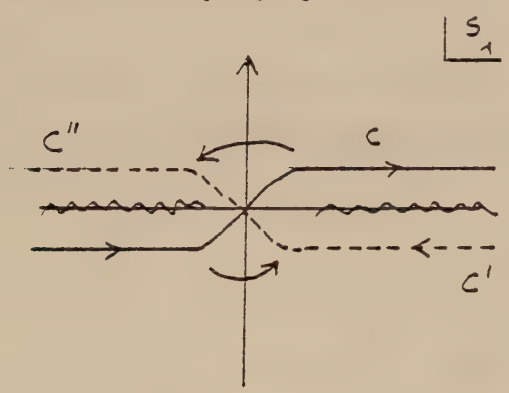


Fig. I .6

Si l'amplitude est telle que l'on peut déformer le contour on peut le rabattre sur lui-même soit à gauche, soit à droite ce qui donne C'. On a une intégrale sur les singularités de la voie s_1, si on intégrait sur C'' ce serait celles de la voie u_1.

On peut donc écrire :

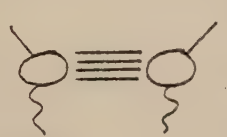
$$N_1(t, t_1, t_2) = \int_0^\infty ds_1 \, \mathrm{Disc} \left\{ I(s_1, t_1, t_1, t_2) \right\}_{s_1} \qquad (I.13)$$

Dans ces conditions, on constate aisément qu'en l'absence de singularités dans la voie s_1 ou dans la voie u_1, l'intégrale donnera une contribution nulle et il n'y aura pas de coupure de Regge dans l'amplitude physique. Il est donc essentiel pour avoir des coupures qu'il y ait des singularités à la fois en s_1 et u_1, ce qui implique que s_1 et u_1 doivent être non exotiques ensemble. La discussion faite sur N_1 se transpose aisément à N_2.

On obtient des règles de sélection pour les coupures de Regge en choisissant certaines formes pour I et J, nous reverrons ces règles de sélection à l'aide des diagrammes de dualité dans le chapitre V.

D'après ce que nous avons vu dans le chapitre II, l'intégrale sur la partie imaginaire d'une amplitude est égale au résidu du premier pôle fixe qui apparaît à $\ell = -1$ dans le cas sans spin. Le vertex de Gribov est donc la contribution de ce pôle fixe dans la réaction Regge + particule $\longrightarrow$ Regge + particule. Ainsi on se trouve ramené à calculer l'amplitude (I .12) qui est proche de la forme habituelle d'une coupure de Regge dans laquelle interviennent les résidus des pôles fixes non sens des réactions $a + R_1 \rightarrow b + \bar{R}_2$ et $c + \bar{R}_1 \rightarrow d + R_2$. A partir des réactions inclusives, on peut espérer obtenir une idée de ces résidus. Des règles de somme ont été obtenues, qui relient ces résidus à des quantités physiques. [6]

N_1 et N_2 sont des intégrales sur les parties imaginaires des amplitudes particule Regge $\rightarrow$ particule Regge. Par exemple pour des Reggeons identiques, on peut relier cette intégrale à la section efficace totale, si on se place à $t = 0$. [7]

$N = \int \Sigma$  $= \int Im$

C'est une application du théorème optique généralisé de Mueller.

Dans le cadre où nous nous sommes placés - technique des graphes de Feynman - pour étudier les échanges de Regge, on peut se demander quels types de graphes donnent des singularités en s et en u.

Mandelstam a montré que les graphes non plans remplissent cette condition et donc le graphe le plus simple qui donne une coupure est le suivant (fig.I .7,a). [5]

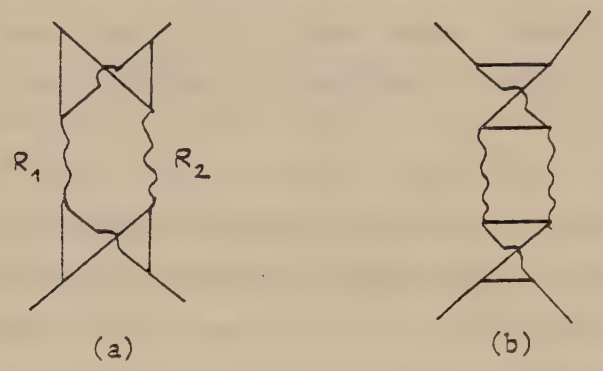


Fig. I .7

Le graphe de la figure I .7,b par contre est réductible à un graphe plan, il ne donnera pas de singularités en s et en u.

Il faut noter le parallèle entre les graphes de Feynman non plans et les "twists" sur les lignes de quarks que l'on met dans les diagrammes de dualité pour relier des échanges successifs. (voir chap. II.)

Cette non planarité est reliée à la présence d'un terme doublement singulier en s et u. Pour la représentation de Mandelstam cela veut dire que la troisième fonction spectrale double $\rho_{su} \neq 0$; dans un modèle de Vénéziano on aura un terme V (s,u).

On peut généraliser la discussion à des échanges d'ordre supérieur. Par exemple pour un échange de 3 Reggeons, le graphe à considérer comme Vertex de Gribov est le suivant :

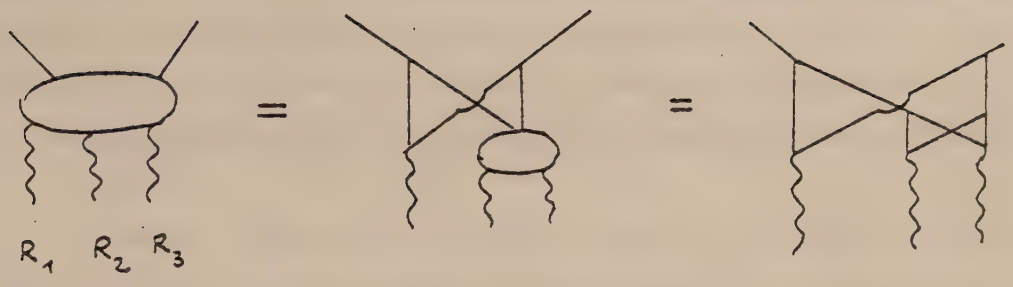


Fig. I .8

Cette règle permet de penser que tous les échanges successifs doivent être séparés par un terme singulier s - u ; dans le chapitre II nous verrons l'application de ceci sur les diagrammes de dualité.

I .3. Singularités du type AFS [5]

Nous nous intéressons à un graphe simplifié où les vertex de Gribov n'ont qu'une particule intermédiaire :

Fig. I .9

Le graphe complet à calculer est le suivant. (fig.I .10)

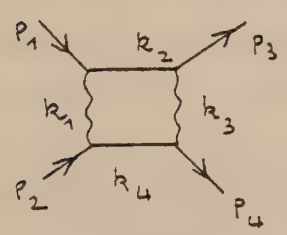


Fig. I .10

k_1 sont les moments internes à la boîte, k_2 et k_4 sont les impulsions de vraies particules. Les vertex de Gribov contiennent un propagateur et les fonctions résidus de Regge.

L'expression de ce graphe est la suivante :

$$A(s,t) = d^4k_1 \int \frac{R_1(k_1^2, k_2^2, k_4^2, s)\, R_2(k_3^2, k_4^2, k_2^2, s)}{\left[k_2^2 - m^2 + i\varepsilon\right] \quad \left[k_4^2 - m^2 + i\varepsilon\right]}$$

(I .14)

R_1 et R_2 sont les amplitudes de Regge associées à l'échange de pôles de trajectoires $\alpha_1(k_1^2)$ et $\alpha_2(k_3^2)$, amplitudes extrapolées hors de leur couche de masse, comme le montre la dépendance dans les moments externes.

Nous faisons le changement de variables suivant :

$$d^4k_1 \rightarrow dk_1^2 \; dk_2^2 \; dk_3^2 \; dk_4^2$$

Le jacobien de la transformation s'écrit :

$$J = \frac{\Theta(D)}{\sqrt{D}} \quad \text{où} \quad D = -16 \, \det \left[2k_i k_j \right] \qquad \text{(I .15)}$$

expression qui pour s grand devient :

$$D \underset{s \to \infty}{\sim} 16 s^2 \left[4 t k_3^2 - \left(k_1^2 - t - k_3^2 \right)^2 \right]$$

Nous posons $k_1^2 = t_1$, $k_3^2 = t_2$ et

$$\tau(t,t_1,t_2) = \frac{\Theta\left(-t^2 - t_1^2 - t_2^2 + 2tt_1 + 2tt_2 + 2t_1t_2\right)}{\left(-t^2 - t_1^2 - t_2^2 + 2tt_1 + 2tt_2 + 2t_1t_2\right)^{1/2}}$$

(I .16)

Ce qui donne :

$$A(s,t) \sim \frac{1}{4s} \int dt_1 dt_2 \; \tau(t,t_1,t_2) \int dk_2^2 dk_4^2 \; \frac{R_1(s,t_1,k_2^2,k_4^2)\, R_2(s,t_2,k_2^2,k)}{\left[k_2^2 - m^2 + i\varepsilon\right]\left[k_4^2 - m^2 + i\varepsilon\right]}$$

(I .17)

L'intégrale sur k_2^2 et k_4^2 contient des singularités : les pôles introduits par les propagateurs et les singularités de R_1 et R_2 dans ces 2 variables. Ces dernières venant de l'extrapolation hors de la couche de masse elles se trouvent dans le 1/2 plan inférieur. [6]

Dans la suite du calcul nous ferons l'hypothèse qu'elles ne viennent pas recouvrir les pôles des propagateurs. Nous séparons l'intégrale en 2 parties : l'une B (s,t) n'a que les pôles des propagateurs, l'autre C (s,t) contient toutes les autres singularités en k_2^2 et k_4^2. Les intégrales sont évaluées respectivement sur les contours C_1 et C_2. Si nous supposons que R_1 et R_2 ont bon comportement à l'infini - du type $k_{2,4}^{-2}$ pour $k_{4,2} \to \infty$ - les intégrales sont bien définies.
Soit alors :

$$A(s,t) \sim B(s,t) + C(s,t) \qquad (I.18)$$

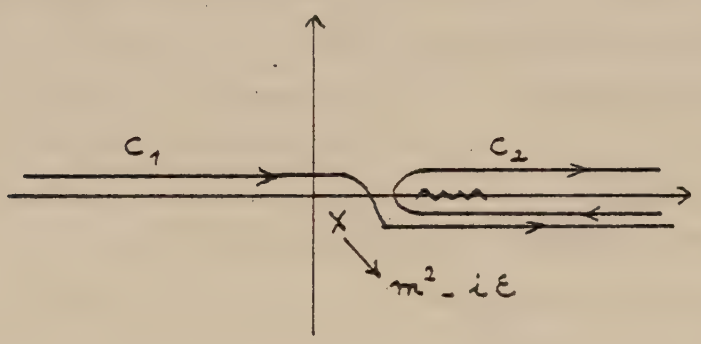


Fig. I .11

Nous l'avons vu au paragraphe précédent, une telle amplitude ne va pas contribuer car les graphes introduits sont plans . Ceci est évident sur la figure I.11, la contribution du pôle est exactement compensée par celle des autres singularités. Ce type de calcul a été d'abord fait par Amati , Fubini et Stanghellini; Mandelstam a prouvé que leur calcul donnait une amplitude nulle et a introduit les graphes non plans , [5] pour éviter leur disparition de l'amplitude physique.

Toutefois nous nous intéresserons à l'intégrale sur le contour C_1 qui donne un comportement de coupures de Regge. L'intérêt de ce terme n'apparaissant pas dans l'amplitude physique n'est pas acadé-

mique. En effet considérons un graphe non plar général qui donne des vertex de Gribov non nuls :

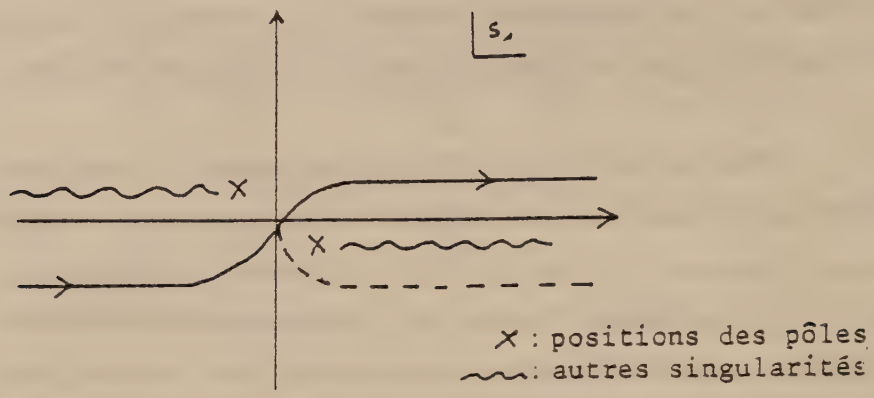


Fig. I .12

Si on peut déformer le contour et si on néglige les contributions autres que celle du pôle, on se trouve ramené au calcul de B (s,t).
C'est précisément ce type de calcul que l'on fait dans les modèles d'absorption, où dans les vertex de Gribov on ne met qu'une particule intermédiaire . On tient compte des singularités négligées en ajustant la force du terme de coupure.

- Calcul de B (s,t)

L'application du théorème des résidus nous donne :

$$B(s,t) = -\pi^2 \int dt_1\, dt_2\, R_1(s,t_1,m^2,m^2)\, R_2(s,t_2,m^2,m^2) \qquad (I.19)$$

ce qui nous ramène au calcul classique de la coupure. L'étude détaillée de ce terme est faite dans l'appendice B.

Nous donnons simplement les résultats pour des trajectoires identiques.

Le comportement asymptotique dépend de la valeur du transfert :

$$t < 4M^2 \qquad B(s,t) \sim D(t)\, S^{2\alpha(t/4) - 1} / \log s$$

$$t > 4M^2 \qquad B(s,t) \sim E(t)\, S^{\alpha(M^2) + \alpha((\sqrt{t} - M)^2) - 1} / \sqrt{\log S} \qquad (I.20)$$

Par une transformation de Mellin , on peut obtenir les singularités du plan complexe l.

En dessous du seuil à 2 particules $t < 4\,M^2$, la singularité dominante est :

$$j = 2\,\alpha\,(t/4) - 1 \qquad \text{(I .21)}$$

Les singularités plus à droite sont dans des feuillets non physiques Quand $t > 4\,M^2$ un autre point de branchement apparaît dans le feuillet physique en passant par le premier point de branchement. Sa position est la suivante :

$$j = \alpha\,(M^2) + \alpha\,((\sqrt{t} - M)^2) - 1 \qquad \text{(I .22)}$$

S'il n'y a pas d'état lié à M , jusqu'à $t = 16\,m^2$ c'est $j = 2\alpha\,(t/4) - 1$ qui domine, ensuite au seuil à 4 particules, on a un nouveau point de branchement à :

$$j = \alpha\,(4\,m^2) + \alpha\,((\sqrt{t} - 2\,m)^2) - 1 \qquad \text{(I .23)}$$

Il faut noter que, le point de branchement à $j = 2\,\alpha\,(t/4) - 1$ est tel qu'il cache le pôle fixe non sens à $\ell = -1$.
En effet pour $t = 4\,m^2$:

$2\,\alpha\,(m^2) - 1 = -1 \Longrightarrow \alpha\,(m^2) = 0$ vérifié pour des particules sans spin

I .4. Interprétation physique. [8]

Nous allons essayer d'interpréter la condition obtenue sur les amplitudes I et J à savoir : les graphes de Feynman associés doivent être non plans ..

Une image physique commode est la théorie de la diffusion du deuteron développée par Glauber. Dans ce cas on pourrait penser que la diffusion est la somme des diffusions sur le proton et le neutron

séparément. La section efficace obtenue par simple addition de ces 2 termes, est trop importante, comparée à la valeur expérimentale; il faut donc introduire un terme correctif - le terme est une double diffusion sur les composants du deuteron. Dans le cas de la diffusion πd, la figure IV.13 montre les termes de diffusion simple .

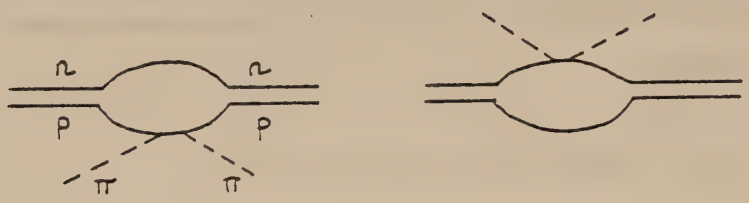


Fig. I .13

Pour les termes de double diffusion (figure I .14) à haute énergie, c'est à dire avec un pion rapide le processus (a) est peu probable, mais (b) peut intervenir :

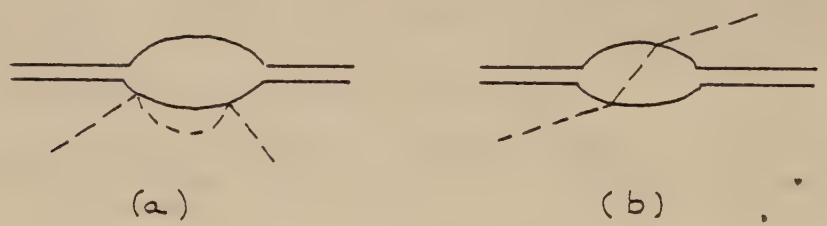
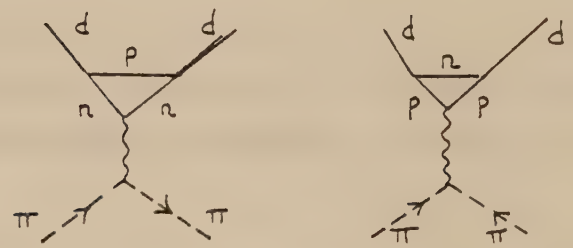


Fig. I .14

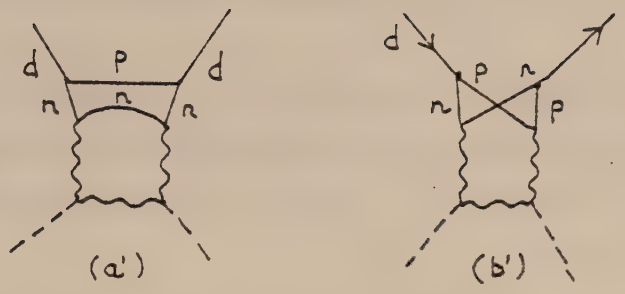
On représente ces processus en graphes de Feynman :



Termes de diffusion simple

Fig. I .15

De même :



Termes de double diffusion

Fig. I .16

On constate que le terme de double diffusion, contribuant à haute énergie, se représente comme un graphe non plan (Au vertex du bas il faut considérer la particule intermédiaire comme composite). Le graphe a' ne va pas contribuer à une coupure.

Ainsi la diffusion π d se représente dans la théorie de Glauber par :

$$A_{\pi d} = G(t)\left[A_{\pi p} + A_{\pi n}\right] + \int G(k^2)\, A_{\pi p}\, A_{\pi n}\, d^2k$$

G (t) est le facteur de forme du deuteron.

On a donc la somme des amplitudes de diffusion sur chaque nucléon individuel et d'un terme de correction, traduisant le fait qu'une partie du temps un nucléon masque l'autre au pion. C'est cette correction qu'on appelle correction d'ombre [shadow scattering] qui correspond à un terme de coupure. Si nous transposons au cas de la diffusion des hadrons, on peut relier la présence de coupures de Regge à la structure des particules qui se diffusent ; pour obtenir un terme de double diffusion, celles-ci se cassent et se reforment ensuite. Cette notion de structure composite des hadrons est pleinement exploitée dans le "dual Reggeon calculus" de Lovelace. Nous pouvons remarquer que cet aspect est

utilisé dans les diagrammes de dualité avec quarks. Nous verrons au chapitre suivant que dans une coupure de Regge, les deux Reggeons échangés sont branchés sur des lignes de quarks différentes. Le mécanisme peut se résumer de la façon suivante : chaque hadron se décompose en quarks (ou en partons selon le modèle) ; ces constituants échangent des Reggeons puis se recombinent en hadrons. L'image physique ainsi obtenue pour les coupures de Regge à l'aide de la diffusion sur des noyaux peut se résumer par les équivalences suivantes :

hadron	$\sim$	noyau
quark ou parton	$\sim$	nucléon
résidu du pôle fixe	$\sim$	fonction d'onde nucléaire
propagateur de Reggeon	$\sim$	amplitude de diffusion nucléon-projectile

APPENDICE A

Singularités des intégrales de fonctions analytiques [2]

Nous donnons ici brièvement les mécanismes qui créent des singularités dans des intégrales de fonctions analytiques. Nous nous intéressons aux intégrales simples.

Soit : $$\int_C F(z,\zeta)\, d\zeta = f(z)$$

$F(z,\zeta)$ est définie sur le contour C et ses singularités sont en $\zeta_i(z)$ points du plan extérieur au contour C.
Quand z varie, ces singularités parcourent des chemins $\zeta_i(z)$ donc $f(z)$ sera bien définie tant que le contour pourra être déformé pour éviter ces singularités.

Dans le cas contraire, l'intégrale sera singulière, on a les 3 possibilités suivantes :

a) Une singularité passe par une extrémité de C c'est une singularité de point final (end point singularity)

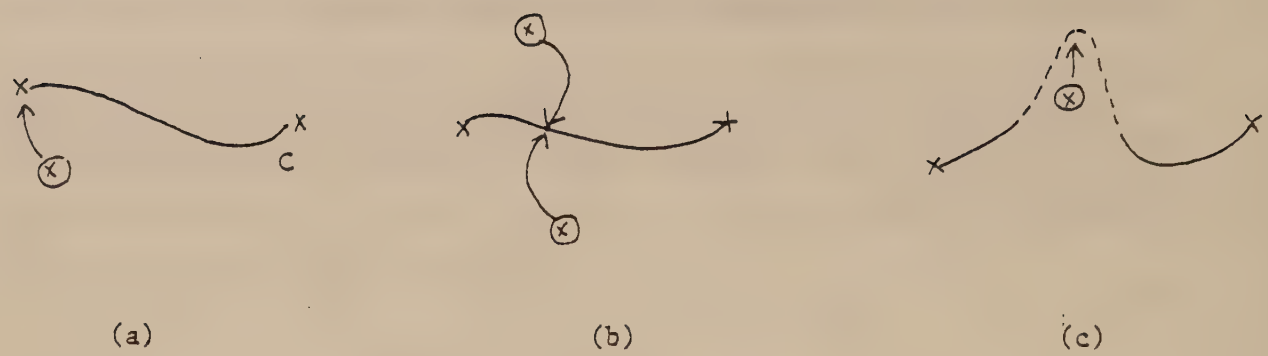
(a)

b) On a 2 singularités de part et d'autre du contour qui viennent coincider en un point $z = z_0$. Le contour est alors pincé par ces 2 singularités : on a une singularité de pincement (pinching singularity)

(b)

c) Une ou plusieurs singularités partent à l'infini et déforment le contour, on peut avoir ainsi une singularité. Toutefois si l'intégrale a un bon comportement à l'infini, le problème disparait.

(c)



(a) (b) (c)

APPENDICE B

Calcul de la "convolution" et du comportement asymptotique des amplitudes.

Considérons la formule (I .19) dans laquelle on explicite les amplitudes de Regge sous la forme (I .4) avec la même trajectoire :

$$B(s,t) \sim -\pi^2 \int dt_1\, dt_2 \quad \tau(t_1,t_2,t)\, \gamma(t_1)\, \gamma(t_2)\, \Gamma(1-\alpha(t_1)) \times$$
$$\times\ \Gamma(1-\alpha(t_2))\, S^{\alpha(t_1)+\alpha(t_2)-1} \qquad (B.1)$$

On fait le changement de variables suivant :

$$x = t_2 - \frac{1}{4t}(t_1 - t - t_2)^2$$

$$r = \frac{t_1 - t - t_2}{2\sqrt{-t}} - \frac{1}{2}\sqrt{-t}$$

C'est à dire :

$$t_1 = x - \left[r + \frac{1}{2}\sqrt{-t} \right] = x - X$$

$$t_2 = x - \left[r - \frac{1}{2}\sqrt{-t} \right] = x - Y$$

Il vient alors :

$$B(s,t) \simeq -\pi^2 \int_{-\infty}^{0} \frac{dx}{\sqrt{-x}} \int_{-\infty}^{+\infty} \gamma(x-X)\, \gamma(x-Y)\, \Gamma(1-\alpha(t_1))\, \Gamma(1-\alpha(t_2)) \times$$
$$\times\ S^{\alpha(t_1)+\alpha(t_2)-1} \qquad (B.2)$$

Considérons tout d'abord la structure analytique dans la variable r. Les singularités viennent de $\gamma(u)$, $\alpha(u)$ et $\sin \pi\alpha(u)$. Si $\sqrt{u_n}$ sont les masses des états liés, des résonances apparaissent sur la trajectoire $\alpha(u)$ et les positions des seuils alors les singularités en r

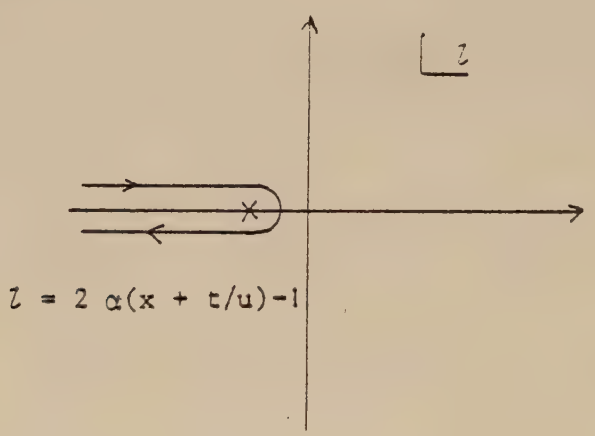
sont situées aux points r_n tels que :

$$r_n = \pm \frac{1}{2} \sqrt{-t} \pm i \sqrt{u_n - x} \qquad (B.3)$$

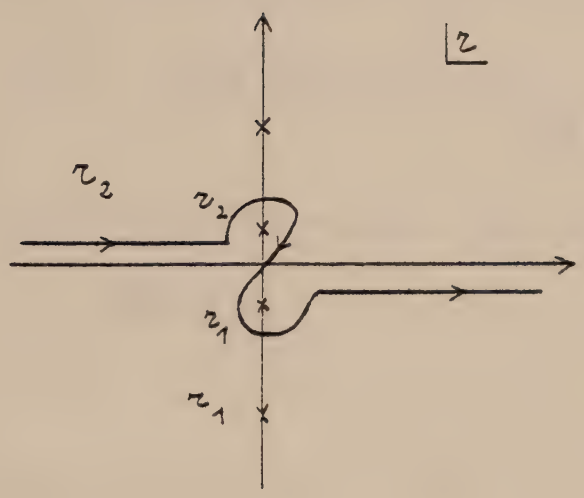
En définissant $\zeta = \alpha (x - X) + \alpha (x - Y) - 1$ on peut écrire :

$$B(s,t) \sim -\pi^2 \int_{-\infty}^{0} \frac{dx}{\sqrt{-x}} \int_{C_\zeta} \frac{d\zeta}{d\zeta/dr} B(\zeta,t,x)\, S^\zeta \qquad (B.4)$$

Le contour C_ζ est un contour dans le plan ζ qui entoure les singularités de l'intégrale. [Pour $-\infty < r < +\infty$, ζ va de $-\infty$ à $-\infty$].
Pour $u < 0$ on suppose les fonctions $\alpha(u)$ analytiques réelles, la seule singularité sur l'axe vient de $(d\zeta/dr)_{r=0} = 0$.
On a donc un point singulier à $\zeta = 2\alpha(x + t/u) - 1$ le contour minimum est le suivant :



Pour prolonger à $t > 0$, il faut repasser dans le plan r où on voit bien apparaître les singularités localisées par (B.3). Elles sont sur l'axe imaginaire à $t = 0$ et s'y déplacent pour $t > 0$.
Si M^2 est la matérialisation la plus basse sur $\alpha(u)$, pour $t = 4(M^2 - x)$ 2 singularités conjuguées r_1 et r_2 viennent au point 0 et pour $t > 4(M^2 - x)$ repartent dans des directions opposées : on a le contour de la figure B.2. Pour $t = 4(M^2 - x)$ on a une singularité de pincement [cf. Appendice A].



(B.2)

- Comportement asymptotique

Le comportement asymptotique dépend des singularités en r et de la limite supérieure de l'intégration sur x soit $x = 0$.
Pour $t > 4M^2$ le comportement est controlé par la singularité à $r = 0$ et $x = 0$ soit $j = 2\alpha(t/u) - 1$.
Pour $t < 4M^2$ c'est par contre $r = \pm i\,(1/2\sqrt{t} - \sqrt{M^2 - x})$ et $x = 0$ c'est à dire $l = \alpha(M^2) + \alpha((\sqrt{t} - M)^2) - 1$ qui donne le comportement asymptotique.
En développant les trajectoires au voisinage des points indiqués et en négligeant les variations de fonctions lentement variables on obtient :

Pour $t < 4M^2$

$$B(s,t) \sim -\pi^3 \left[\Gamma(1 - \alpha(t/4))\,\gamma(t/4)\right]^2 \frac{S^{2\alpha(t/4) - 1}}{2\alpha'(t/4)\log S}$$

Et $t > 4M^2$

$$B(s,t) \sim 2\pi^{s/2} \frac{\gamma((\sqrt{t} - M)^2)\,\gamma(M^2)\,\Gamma(1 - \alpha(M^2))\,\Gamma(1 - \alpha((\sqrt{t} - M)^2))}{M\,\alpha'(M^2)\,\sqrt{\frac{t}{M}\alpha'((\sqrt{t} - M)^2)}}$$

$$\frac{S^{\alpha(M^2) + \alpha((\sqrt{t} - M)^2) - 1}}{\sqrt{\log S}} \qquad (B.5)$$

Dans le cas général où les trajectoires ne sont pas les mêmes pour

$t < 4\ M^2$ la position de la singularité est :

$j = \alpha_1 (u) + \alpha_2 (\sqrt{t} - \sqrt{u})^2) - 1$ où u est une solution de $d^2 /dr = 0$ soit :

$$\alpha_1' (u) - \alpha_2' ((\sqrt{1/2} - \sqrt{u})^2) (\sqrt{t} - \sqrt{u}) / \sqrt{u} = 0$$

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CHAPITRE II

DUALITY DIAGRAMS AND REGGE CUTS IN MESON-BARYON SCATTERING – DESCRIPTION OF π-N AND K-N SCATTERING

by

G. GIRARDI, R. LACAZE, R. PESCHANSKI

ABSTRACT

We propose a reggeized absorption model including lower lying Regge cuts and satisfying duality requirements. It is shown that the introduction of these singularities, especially the Reggeon-Pomeron-Reggeon (R-P-R) cuts, provides a solution to many puzzles in meson-baryon scattering. Theoretical evidences for the cuts, using duality diagrams and s-channel unitarity are given. These theoretical constraints reduce strongly the flexibility of the model and give it a good predictive power. With only 12 parameters, we obtain an excellent description of observables in π-nucleon and K-nucleon forward scattering in a large energy range ($5 \lesssim P_{lab} \lesssim 30$ GeV/c). This success provides a better understanding of Regge cuts in a duality framework which can be interpreted as a phenomenological approach to a perturbative dual field theory.

INTRODUCTION

Until recently, it was thought that reggeized absorption models could be successful in describing two-body or quasi two-body reactions. With a few Regge-pole exchanges (R),and the Pomeranchuck singularity (P) describing the high energy limit of elastic scattering, one achieved a good understanding of almost all data, provided the Regge-pole exchanges were corrected by absorption through the convolution R * P . But a better knowledge of the actual amplitudes due to more precise or new results ($\pi^- p \to \pi^o n$ polarization) rules out the Reggeized absorption models, in which the Pomeron has as simple a structure as a normal pole.

Among different attempts, one can take the attitude that all that is needed to explain the data are corrections to the reggeized absorption models. There are three possibilities : either change the input poles (R or P), or change the absorption corrections (R * P), or add new contributions, that is lower lying singularities. Our model belongs to the third class and its typical features are the economy in the number of free parameters and the strong and direct support it takes from duality. Duality makes it possible to consider secondary Regge cuts, with multiple exchange of Regge poles and Pomerons, in such a way that their parametrization, or their existence itself in the studied channels, are submitted to strong selection rules.

Our main result is that a special kind of secondary cuts - R*P*R - , and it alone, is able to solve the puzzles of absorption models. It has the following properties :

- it gives a central (non peripheral) real part in the negative signature amplitude.
- It increases the absorption in the imaginary part,
- It gives a violation of exchange degeneracy, that is different contributions in the imaginary parts of negative and positive signature amplitudes.

In fact, our model provides an answer to one of the questions arising from a crucial problem of duality : how to conciliate duality and Regge cuts ? The links that exist between the phenomenological and more theoretical puzzles of strong interactions are discussed in the motivations of our model (section 1), which express the "philosophy" of our approach.

Duality, in this framework, is a perturbative theory where the "Born term" is given by the Regge pole duality diagrams (section 2). From this term we build the expansion of the high energy amplitude and find (already known and new) selection rules for Regge cuts in meson baryon scattering (section 2 .1). A perturbative view of n particle production amplitudes makes it possible to use unitarity and to go deeper into the physical understanding of the nature and properties of the new proposed singularities - the R*P*R cuts - (section 2.2).

The description of πN and KN reactions, which is proposed as a practical application of the model (section 3), uses as much as possible the theoretical framework. In particular the parametrization of the s-channel helicity amplitudes (section 3.1) involves only 12 parameters for 8 studied reactions ($\pi^- p \to \pi^- p$, $\pi^+ p \to \pi^+ p$, $\pi^- p \to \pi^0 n$, $\pi^- p \to \eta n$, $K^+ p \to K^+ p$, $K^- p \to K^- p$, $K^- p \to \bar{K}^0 n$, $K^+ n \to K^0 p$). A fit, performed on forward elastic pion-nucleon data, leads to a very good description of available data of all the reactions considered in the range $-2.5\ \text{GeV/c} \leq P_{lab} \leq 40\ \text{GeV/c}$, and $0 \leq -t \leq 1\ (\text{GeV/c})^2$ (section 3.2).

At last (section 4) the conclusions are drawn, and some possible experimental tests and some theoretical open problems are discussed.

1 - MOTIVATIONS OF THE MODEL

The introduction of reggeized absorption models (and thus of Regge cuts), the invention of the duality concept (with the discovery of a new symmetry property of Regge trajectories and residues, namely exchange degeneracy) have been very promising developments of Regge phenomenology. But during the recent past this field of research has known a sort of a standstill.

The difficulties concerning its development are both of theoretical and phenomenological nature. On pure theoretical grounds one must recognize that when they were introduced, reggeized absorption models were lacking a real justification. Several questions were not solved : reggeization does not already take into account some "absorption" ? Thus, is it not double counting to include absorptive cuts ? Moreover a Regge cut model suggests a perturbation expansion whose Born term is the Regge pole exchange. Now, in field theoretical models, reggeization comes out from summing infinite series of high order diagrams[1]. Therefore perturbative Regge cuts do not seem theoretically justified. Nevertheless important theoretical progress has been made in understanding Regge cuts. The Reggeon calculus, mainly developed by Gribov[2] projects some light on Regge cuts;unfortunately the discontinuity of the cut is given by an integral over an off-shell Reggeon-Particle amplitude, about which physical information is lacking. On the other hand the derivation of duality, originally based on pure phenomenological considerations, suffers from lack of theoretical justification. For instance, duality was entirely derived under the assumption that Regge cuts can be neglected, which is in contradiction with phenomenological observations (πN exchange polarization, π exchange).

Very severe phenomenological difficulties appeared in a more detailed analysis, which discouraged developing both reggeized absorption models and dual models. For instance all simple absorption models failed in predicting the πN charge exchange polarization at large values of $|t|$ ($|t| \gtrsim 0.2(\mathrm{GeV/c})^2$). Moreover precise amplitude analyses which are now possible from complete set of measurements have shown that exchange degeneracy, which works well for the materialization of resonances, is definitely broken in the space-like region (KN and $\bar{K}N$ charge exchange reactions, hypercharge exchange reactions ...).

Despite these troubles, and on the basis of the impressive qualitative successes of absorption and duality, we do not think that one should give up investigations in these two fields. Moreover it appears that it is necessary and enough to add to the exchange degenerate poles (eventually corrected for absorption)

a well defined contribution, more "central" (important in lower partial waves) and possibly with smaller intercept than the leading Regge singularity[3]. So that, we claim that solutions to the difficulties of the reggeized absorption models can be found in duality and that troubles in duality requirements can be explained in terms of Regge cuts. In fact it has been already observed that absorption implies some effects which are of a "dual" type. The Pomeron plays a special role in absorption as well as in duality. Absorption depends crucially on s-channel helicity quantum numbers, so, in absorption models the high energy behaviour of amplitudes depends on s-channel quantum numbers, which is a dual property. Actually people have proposed simple minded geometrical parametrization which they ambitiously call "dual absorptive models". On the other hand duality has not remained a purely phenomenological construction. It has been proposed that the narrow resonance dual amplitudes (generalized to multiparticle production amplitudes) can be considered as the Born term of a new type of field theory, the dual field theory (D.F.T.). Indeed this ambitious program (the outcome of which would be unitary, crossing symmetric dual amplitudes !) encounters very great difficulties related to the renormalization of multiloop diagrams ; however, for phenomenelogy (which, after all, does not have to wait for the solution of all theoretical problems) very important informations can be extracted from D.F.T. First of all, in D.F.T., the Born term is already Regge pole behaved. Planar loops are expected to affect neither the pole nature of Regge singularities nor exchange degeneracy. On the contrary, non-planar loops exhibit Regge cuts and exchange degeneracy breaking. More precisely it has been shown that one can derive from D.F.T. a dual Reggeon calculus[4], with selection rules, models for the Reggeon-particle amplitudes and thus the Gribov vertices. Moreover the Pomeron appears as an output of D.F.T., with some understanding of the Pomeron-background duality. It thus appears that D.F.T. is precisely the theoretical framework one needs to derive reggeized absorption models including Regge cuts responsible of the breaking of exchange degeneracy.

It is now possible to formulate precisely the philosophy which is at the origin of the model we propose :

i) Our model is a high energy perturbative expansion based on D.F.T. requirements

ii) The "Born" term of this expansion is taken as pure Regge pole exchange, satisfying strong exchange degeneracy (actually our Regge pole contribution is not necessarily the same as that given by the high energy behaviour of narrow resonance dual amplitudes, since planar loops can change both the trajectory and the residues. That is why the scale parameter ν_o is not fixed at $\frac{1}{\alpha'}$).

iii) In the absence of any further information, the Pomeron is parametrized as a simple pole with intercept 1.

iv) All Regge cuts lying one half unit below the dominant (non Pomeron) singularity are taken into account : that is Regge-Pomeron (or absorptive) cuts, Regge-Regge cuts and Regge-Pomeron-Regge cuts. Indeed a complete theory should include all Regge cuts. But such a theory is for the moment out of reach since already the Regge pole contribution is parametrized to order $1/s$.

v) Regge cuts are parametrized in such a way as to satisfy the selection rules implied by non planar loop dual diagrams.

vi) In the absence of a complete dual Reggeon calculus, Gribov vertices are parametrized in terms of multiplicative real constants. Qualitative arguments, mainly based on multiproduction contributions to the s-channel unitarity condition are used for guessing the sign and eventually the strength of these constants.

In a previous paper[5] we have used such a philosophy to introduce Regge-Pomeron-Regge (R*P*R) cuts, in particular in πN charge exchange. We show that the apparently catastrophic failure of reggeized absorption models to describe the polarization is actually a small effect, in the sense that it can be corrected through the use of low lying Regge singularities. Provided that the sign of its contribution is conveniently chosen, the R*P*R cut is enough to produce the desired effect. In the present paper we generalize previous arguments to include an extensive study of all high energy data for πN, KN and $\bar{K}N$ scattering. Moreover we present an argument based on s-channel unitarity justifying the correct sign for the R*P*R contribution.

Before going into the details of the derivation of the model, it is useful to insist on the features which are in favor of our attitude as compared to that which is at the basis of other models.

The use of duality constraints restricts as much as possible the number of free parameters. So our model constitutes a serious confrontation of Regge cut theories with experiment. This is a very important aspect : Regge cuts are intimately related to the compositeness of hadrons[4], so, any evidence of the presence of Regge cuts in the actual scattering amplitude is indeed welcome. One must recognize that more than 8 years after the introduction of Regge cuts through absorption models, the question of the existence itself of Regge cuts

is still open (alternative models using complex poles[6], ρ' pole[7]..., continue to be proposed). In this respect our model can be considered as a new type of test of Regge cut theories. The economy of the parametrization has to be compared with that of the so-called i factor model[8] which provides rather nice fits but appears to be somewhat ad hoc.

In some sense our model can also be considered as a phenomenological exploration of the validity of D.F.T. After all, difficulties in D.F.T. are so large that one may ask the question whether it is useful to expand so many efforts in a field which could turn out to be a pure academic exercise. Our phenomenological point of view could provide some encouragement for such investigations. It is opposite to the one underlying the attempts to improve the so-called SCRAM model[9], in which one tries to make the Pomeron responsible of all the troubles of the original model but with no obvious connection with duality.

2 - DERIVATION OF THE MODEL

The ideas which guide us, for the derivation of the model are the following. The comparison between the actual amplitudes and those of reggeized absorption models[3] exhibits the main properties we expect from the secondary contributions we introduce.

- They should exist in exotic s and u channels ; this is an important feature since many troubles appear in such reactions (K^+N and $\bar{K}N$ scattering, hypercharge exchange).

- They should contribute to negative signature amplitudes and in particular to the real parts of low partial waves (central contribution).

- They should show a violation of exchange degeneracy. The method we use to introduce the lower-lying singularities is derived from a perturbative approach based on duality diagrams. In two body meson-baryon scattering one obtains directly selection rules for the amplitudes by describing each t-channel Regge singularity by a simple duality diagram exhibiting that same singularity together with the dual singularities in the other channels. This is done in details in section 2.1. But, many diagrams exhibit the same singularity, among which are all those that are obtained from the simplest one by adding any number of loops. In particular one may show that the Pomeranchuck singularity is built from such a sum of graphs[10], and use s channel unitarity to compute

it in the framework of the multiperipheral model. The question of cuts containing the Pomeron - R*P and R*P*R cuts - can be discussed in the same way, as is shown in section 2.2, provided one modifies the pure multiperipheral amplitude in the unitarity relation.

2.1. Duality diagrams in two body meson baryon scattering

Our starting point is the perturbative expansion of the amplitude in terms of Regge poles and their iterations.

$$M = P + R + P*R + R*R + R*R*P + R*P*R + \ldots \qquad (1)$$

The poles (resonances or Regge poles) correspond to non exotic intermediate quark states in the diagram ($q\bar{q}$ for mesons and qqq for baryons).

The Regge cuts (fig. 2.1.a) are built from the prescriptions of Mandelstam and Gribov[2],[11] which state that the amplitudes for the subprocesses I and J (fig. 2.1.b,c), the so-called "Gribov vertices", have singularities in both their s and u channels. Concerning duality diagrams, this prescription leads us to choose as Gribov vertices the diagrams having poles in these channels, the (s,u) term (fig. 2.1.d). More complicated diagrams of dual field theory can be used, but the diagrams we choose have the advantage of being simple and computed in the dual Reggeon calculus[4]. We will discuss this approximation in the conclusions of our study (section 4).

Pole exchange : P,R

Regge pole exchanges (R) are represented by the duality diagrams given in figure 2.2 (a,b). The main consequence of using duality diagrams for Regge poles is strong exchange degeneracy of SU(3) nonets of mesons with opposite signature. The (u,t) diagram (fig. 2.2.a), has poles in the t and u channels, but no poles (resonances) in the s-channel. It corresponds to the real combination of two exchange degenerate poles (denoted by $\mathbb{1}$). The (s,t) diagram, obtained by s-u crossing from the previous one (fig. 2.2.b) has a rotating phase and s channel resonances (we denote it by $e^{-i\pi\alpha}$).

The vacuum exchange (P) is represented by the one loop, non planar diagrams of fig. 2.2.(c,d). These diagrams are in agreement with the required properties of the Pomeranchuck singularity, namely no exchange of quantum numbers in the t channel and no resonances in the s and u channels. In D.F.T. the

background to which the vacuum exchange is dual, is made of Regge cuts. Vacuum exchange being s-u symmetric, both diagrams c and d represent the same contribution P.

One notices from the diagrams of figure (2 .2) that poles exist in a channel whenever a quark joins one particle to another in that same channel. We shall call this a "quark line annihilation", since it can be interpreted as a $q\bar{q}$ annihilation. The diagrams (a) or (b) have quark-line annihilations in two channels, (u,t) or (s,t), whereas the diagrams (c) and (d) have two quark-line annihilations in the same t channel. In the same way, the diagrams of fig. .1.d have annihilations in both the s and u channels.

Regge cuts : P*R, R*R, R*R*P, R*P*R

Starting with figure (2 .1), we obtain the diagrams corresponding to Regge cuts by connecting the pole diagrams by (s-u) terms (fig. 2 .1.d). This is equivalent to twisting the quark lines of the intermediate particles. We apply it as a systematic rule, in order to build the different Regge cuts of expansion (1) represented by the diagrams of figure (2.3) : P*R (3-a), R*R(3-b), R*R*P(3-c), R*P*R(3-d).

Though they are not pole diagrams, the duality diagrams of figure (2 .3) can also be classified in terms of quark line annihilations by looking for the path of each quark in the diagram. The results of this classification are given in table 1. From this table can be read off easily selection rules for the diagrams and consequently for the Regge cuts they represent : a Regge cut diagram exists whenever the corresponding pole diagram (with the same quark line annihilations) can be drawn. This rule makes it possible to extend the idea of "exotism" to Regge cuts, since it corresponds to forbid not only a particular pole diagram but also the corresponding cut diagrams of table I. Let us now consider in detail each singularity and point out the phenomenological implications of the selection rule.

Absorptive cuts (R*P) follow the same selection rules than the associated Regge poles. In particular both Regge poles and their absorptive corrections are forbidden in exotic channels. In fact the absorption corrections to exchange degenerate poles give only a small violation of exchange degeneracy, due to the small real part of the Pomeron.

Regge-Regge cuts (R*R) (fig. 2 .3.b)

From table I, one can show three important properties of R*R cuts i) they don't exist in reactions for which the s or u channel is exotic, ii) R*R cuts tend to cancel in negative signature amplitudes, iii) R*R cuts do not give a violation of exchange degeneracy.

i) The selection rules of R*R cuts are those of an (s,u) term which is not present for exotic s or u channels. The best example is KN and $\bar{K}N$ reactions, as already stated by Finkelstein[12]. However the rule is more general and can apply to non exotic reactions, such as $\pi\Xi$ elastic scattering, where the (s,u) term cannot be drawn.

ii) In reactions for which R*R cuts are present (e.g. πN scattering), they have the form "$1*e^{-i\pi\alpha}$". Consider the degenerate Regge poles R^+, R^-. As a convention one may take $R^+ + R^-$ as the combination corresponding to the diagram "1" and $R^+ - R^-$, to the diagram "$e^{-i\pi\alpha}$". For two pairs of degenerate poles $R_1^{\pm}$ and $R_2^{\pm}$, we get in the negative signature amplitude M^- (using the antisymmetric combination of two (s,u) terms of table I)

$$M^- = (R_1^+ - R_1^-)*(R_2^+ + R_2^-) - (R_1^+ + R_1^-)*(R_2^+ - R_2^-) = 2\,(R_1^+ * R_2^- - R_2^+ * R_1^-) \quad . \qquad (2)$$

This gives a partial or total cancellation, as already established by Worden using the Chan-Patton prescriptions[13].

iii) When we consider the dominant Regge trajectory (with $\alpha \simeq 1/2$) the contribution of "$1*e^{-i\pi\alpha}$" is mainly real, and does not help in describing the apparent violation of exchange degeneracy. These various properties show that R*R cuts are not able to solve the phenomenological difficulties of the absorption models.

Regge-Regge-Pomeranchuck cuts (R*R*P) (fig. 2..3.c) . R*R*P cuts can be interpreted as the absorptive corrections to R*R cuts and have the same properties.

Regge-Pomeranchuck-Regge cuts (R*P*R) :(fig. 2 .3.d)

Inserting the Pomeranchuck pole in between the Reggeons give quite different properties from those of the R*R*P cuts : i) R*P*R cuts are present

in exotic channels ii) R*P*R cuts give no cancellation in negative signature amplitudes iii) They give contributions violating exchange degeneracy.

i) In table I, R*P*R cuts correspond to (s,t) or (u,t) diagrams. In an exotic s-channel, say, one can draw the (u,t) term, and then the corresponding R*P*R diagram : " $\mathbb{1}*P*\mathbb{1}$ ". The same is true for an exotic u-channel, with the (s,t) term and the diagram : "$e^{-i\pi\alpha}*P*e^{-i\pi\alpha}$".

ii) The negative signature is obtained by the combination of these two different diagrams. One gets, in the same way as equation (2) :

$$M^- = \left\{ (R_1^+ + R_1^-)*P*(R_2^+ + R_2^-) - (R_1^+ - R_1^-)*P*(R_2^+ - R_2^-) \right\} = \\ = 2(R_1^+ * P * R_2^- + R_1^- * P * R_2^+) \quad . \tag{3}$$

R*P*R cuts give a non vanishing contribution to M^- . As a double convolution of Regge poles, their contribution is central.

iii) The violation of exchange degeneracy is obvious in the exotic channels. The cut "$\mathbb{1}*P*\mathbb{1}$" is namely imaginary, as well as the Pomeranchuck singularity itself.

$$\mathrm{Im}\,(\mathbb{1} * P * \mathbb{1}) = \mathbb{1} \times \mathrm{Im}\,P * \mathbb{1} \; .$$

These properties show that R*P*R cuts make it possible to describes the physical amplitudes in the framework of an absorption model with lower lying Regge cuts satisfying duality requirements. This makes it differ from a previous model[14], for which R*R cuts play the main role.

2.2. Regge cuts from s channel unitarity

Concerning R*P*R cuts, which are the main new singularities we have introduced, arises the same question as for the absorption corrections : what are the physical arguments for introducing such terms and what can be said about their sign and their strength with respect to the other singularities ? Actually, in its present status dual theory of two-body scattering does not answer these questions.

It is tempting to refer to the same physical picture as is relevant for absorption : the reflection of the many body production channels in the two

body scattering by means of unitarity. In the following we propose a generalization of the shadow scattering picture which, in addition to the usual absorption correction provides us with secondary contributions in agreement with dual selection rules.

The Pomeranchuck singularity can be described as the shadow scattering of the many particles channels opened at very high energy, through the unitarity equation

$$\sum_n A^J_{fn} \; A^{J\dagger}_{ni} = \text{Im}\, P^J \, \delta_{fi} \quad . \tag{4}$$

Equation (4) exhibits the t channel partial waves of the Pomeranchuck amplitude (P^J) and multiproduction amplitudes (A^J). Here we have implicitly used the coherence condition of Van Hove[15], that is only diagonal terms (the initial and final states are the same, i ≡ f) which are real positive, give a non vanishing contribution. In the following, we consider that A_n is described by the multiperipheral model (fig.2.4). In order to study, through shadow scattering, the cuts in which the Pomeranchuck appears, we propose a correction to the T_{ni} amplitudes, such that :

$$T_{ni} = A_{ni} + \{R * A\}_{ni} \qquad \text{see fig.2.4.} \tag{5}$$

This form is nothing else than the two first terms of the Regge perturbative expansion of the production amplitudes. This is particularly true in dual perturbative theory where the Pomeranchuck singularity is built by dual multi-Regge diagrams[10] .

Substituting T_{ni} for A_{ni} in the unitarity equation (4) gives new contributions to the imaginary part of the two body amplitudes, with different order in the R interaction, of which we shall give a more explicit form

$$\sum_n T^J_{fn} \, T^{J\dagger}_{ni} = \sum_n \left[A^J_{fn} + (R*A)^J_{fn} \right] \left[A^{J\dagger}_{ni} + (R*A)^{J\dagger}_{ni} \right]$$

$$= \left[\sum_n A^J_{fn} A^{J\dagger}_{ni} \right] + \left[\sum_n (R*A)^J_{fn} \, A^{J\dagger}_{ni} + A^J_{fn} \, (R*A)^{J\dagger}_{ni} \right] + \tag{6}$$

$$+ \left[\sum (R*A)^J_{fn} \, (R*A)^{J\dagger}_{ni} \right] = Y_P + Y_{RP} + Y_{RPR} \quad .$$

Here Y_p is the imaginary part of the Pomeranchuck amplitude and we will show that Y_{RP} and Y_{RPR} are associated respectively with absorption corrections and the RPR cut. Using the convolution formula (see Appendix 1) with one intermediate state, one can calculate Y_{RP} and Y_{RPR}. One finds

$$(R*A)^J_{fn} = \frac{ip}{8\pi\sqrt{s}} R^J_{fh} A^J_{hn}$$
$$(R*A)^{J\dagger}_{ni} = -\frac{ip}{8\pi\sqrt{s}} A^{J\dagger}_{nk} R^{J\dagger}_{ki} \quad . \tag{7}$$

ii) Using equations (4) and (7), Y_{RP} can be interpreted as absorption corrections :

$$Y_{RP} = \frac{ip}{8\pi\sqrt{s}} \left[R^J_{fh} \left(\sum A^J_{hn} A^{J\dagger}_{ni} \right) - \left(\sum_n A^J_{fn} A^{J\dagger}_{nk} \right) R^{J\dagger}_{ki} \right]$$
$$= \frac{ip}{8\pi\sqrt{s}} \left[R^J_{fh} \, \mathrm{Im}\, P^J \delta_{hi} - \mathrm{Im}\, P^J \delta_{fk} R^{J\dagger}_{ki} \right] \quad . \tag{8}$$

With the approximation $\mathrm{Re}P = 0$ we obtain

$$Y_{RP} = \frac{ip}{8\pi\sqrt{s}} \mathrm{Im}\, P^J (R^J_{fi} - R^{J\dagger}_{fi}) = -\frac{2p}{8\pi\sqrt{s}} \mathrm{Im}\, P^J \, \mathrm{Im}\, R^J_{fi}$$
$$\simeq \mathrm{Im}\, (2i(P*R)^J_{fi}) \, .$$

iii) Similarly we rewrite Y_{RPR} :

$$Y_{RPR} = \left| \frac{ip}{8\pi\sqrt{s}} \right|^2 R^J_{fh} \left(\sum A^J_{hn} A^{J\dagger}_{nk} \right) R^{\dagger J}_{ki} \quad . \tag{9}$$

In exotic reactions Y_{RPR} corresponds to the Regge-Pomeranchuck-Regge duality diagram : $\mathbb{1} * P * \mathbb{1}$, since the Regge pole contribution is of (u,t) type in these reactions. Using the definition of the convolutions (equation 7), one can transform (9) to obtain

$$Y_{RPR} = \left(\frac{p}{8\pi\sqrt{s}} \right)^2 \mathbb{1}^J_{fh} \, \mathrm{Im}\, P \, \mathbb{1}^J_{hi} = -\mathrm{Im}\, (\mathbb{1} * P * \mathbb{1})^J_{fi} \quad . \tag{10}$$

In the non exotic reactions, the Regge pole contributions are of (s,t) type, thus not real in the s-channel. So that the preceding trick cannot be used. However s-u crossing relates exotic and non exotic channels, i.e. $1 * P * 1$ and $e^{-i\pi\alpha} * P * e^{-i\pi\alpha}$ contributions. Then, once existing in an exotic channel, the R*P*R cut also exists in non exotic channels, with the form predicted by duality selection rules.

With the proposed correction to T_{ni} (equation 5) we only obtain part of the total two body amplitude, since Im R and some contributions to the R*P*R cuts are lacking. Indeed, we never claimed to get directly a solution of the whole unitarity problem. In particular, we did not pay a detailed attention to the contributions of the two body channels themselves to unitarity.

Our statement is that solving exactly the unitarity equation will not change the singularities which are determined by the inelastic channels, but only give the lacking contributions. The effect of production channels may be well understood when it contributes to background, say in exotic channels. In this precise case, equation 6 can be interpreted as a genuine solution of unitarity since R and R*P do not contribute to the discontinuity of the amplitude, likewise for R*R and R*R*P in agreement with duality prescriptions (see section 2.1). In non exotic reactions where background and resonances occur, the situation is less clear, since ImR itself, is not an output of our calculations. The difficulty for exhibiting exactly the R*P*R term seems related to the peculiar nature of such a contribution which is dual of both resonances and background (fig. 2.5.d).

Our approach is quite similar to Caneschi's who first introduced the same kind of corrections in many body scattering[16]. We differ in the application, since no use is made of unitarity as an iteration procedure. This allows us to stay compatible with duality and the physical picture of our cut corrections. Among the differences, let us note that we do not take care of the elastic part of unitarity, and of $P*A_n$ contributions. In our scheme they would correspond to P*P cuts in the two-body amplitude and only modify the Pomeranchuck singularity.

A new feature of R*P*R cuts appears in equation (10) : their negative sign. We point out that this result is quite interesting when related to the continuous war between the "plus sign" and "minus sign" of Regge cuts. In our derivation, R*P cuts appear with the "plus sign" in agreement with the "eikonale" or "absorption" picture, but R*P*R cuts have the "minus sign", in agreement

with the "multiperipheral" or "rescattering" picture. This behaviour is well justified by the phenomenological analysis of section 3 in pion-nucleon and Kaon-nucleon scattering. Such a feature rules out a general eikonale expansion, even modified to take into account duality, since it would predict always the plus sign.

The strength of R*P and R*P*R cuts can be discussed by taking into account all intermediate states (h) or (k) which can appear in the convolutions (equations 7). Contrary to what is sometimes obtained in other absorption models[9], the introduction of intermediate excited states does not change the strength of absorption in our unitary framework. The summation over these states in equation (8) for Y_{RP} is cancelled by the δ-functions (δ_{hi} and δ_{kf}), which are there because of the coherence condition (4). Only the fundamental states (h = i, k = f), remain, thus giving an absorption strength equal to 1, in agreement with dual Reggeon calculations[4].

The result for the strength of R*P*R cuts is somewhat different. Since taking into account intermediate states involves a double summation on (h) and (k) in equation (9), the coherence condition (4) gives non vanishing contributions for all diagonal terms (h = k). Therefore the strength of R*P*R cuts can be modified, in particular when the couplings of both Reggeons to intermediate states are of the same sign. In this case (for instance for two identical R-exchanges) the intensity of R*P*R cuts can be notably increased. This makes it possible to consider that R*P*R cuts are strongly coupled to the physical channels.

3 - APPLICATION TO πN AND KN SCATTERING

3.1. Parametrization of the amplitudes

We have to parametrize the different contributions of the Regge expansion (1), in agreement with the corresponding duality diagrams for πN and KN scattering.

a) Regge poles (R). As previously mentioned the use of duality diagrams implies the existence of pairs of exchange degenerate Regge poles, with natural parity in πN and KN reactions. The parametrization we choose is in agreement with the high energy limit of a Veneziano formula. The residues are constant, up to an overall scale parameter ν_o ($\nu_o \neq \frac{1}{\alpha'}$)

$$R_i^{\pm} = g^i (\sqrt{-t})^n \left(\frac{\nu}{\nu_o}\right)^{\alpha(t)} \Gamma(1-\alpha(t))(1 \pm e^{-i\pi\alpha(t)}) , \qquad (12)$$

where $\nu = (s-u)/2$, i is the t-channel isospin, and n the total helicity flip in the s-channel.

Writing down the Regge poles and their combinations corresponding to the (u,t) and (s,t) duality diagrams, we define

$$\text{Regge poles : } R_o^+ = P' \quad R_o^- = \omega \quad R_1^+ = A_2 \quad R_1^- = \rho$$

and we obtain for

$$\begin{aligned} &\text{(u,t) diagrams : } r_o = \frac{P'+\omega}{2} \qquad r_1 = \frac{P'-\omega}{2} \\ &\text{(s,t) diagrams : } \varphi_o = \frac{P'-\omega}{2} \qquad \varphi_1 = \frac{A_2-\rho}{2} \end{aligned} \qquad (13)$$

One also defines the inverse formulae

$$P' = r_o+\varphi_o \ , \ \omega = r_o-\varphi_o \ , \ A_2 = r_1+\varphi_1 \ , \ \rho = r_1-\varphi_1 \ .$$

Exchange degeneracy has three important consequences.

i) a unique Regge trajectory,

$$\alpha_\rho(t) = \alpha_{A_2}(t) = \alpha_{P'}(t) = \alpha_\omega(t) = \alpha + \alpha' t \ .$$

ii) exchange degeneracy of intermediate mesons $(\pi,B)(K,Q)$ in Gribov vertices

Since Regge poles have different G-parity in a same degenerate pair, exchange degeneracy cannot give a relation between Regge pole couplings in $\pi N \to \pi N$, where positive G-parity is selected. One has to relate $\pi N \to \pi N$ to the reaction $\pi N \to BN$, where B stands for the Bouddha resonance, the degenerate partner of the pion. In fact, using (s,u) terms to describe the Gribov vertices (fig. 2.1.a) also implies degenerate intermediate particles. Therefore we consider the degenerate pairs of mesons with unnatural parity $(\pi,B),(K,Q)$.

iii) exchange degeneracy and SU_3 relations

The couplings of Regge poles to external particles are related by exchange degeneracy and SU_3 invariance. One obtains :

$$(\rho, A_2) \quad \begin{cases} g_{\rho N\bar{N}} = g_{A_2 N\bar{N}} \\ g_{\rho\pi\pi} = g_{A_2\pi B} = 2g_{\rho K\bar{K}} = 2g_{\rho Q\bar{K}} = 2g_{A_2 K\bar{K}} = 2g_{A_2 Q\bar{K}} \end{cases} \tag{14}$$

$$(\omega, P') \quad \begin{cases} g_{P' N\bar{N}} = g_{\omega N\bar{N}} \\ g_{P'\pi\pi} = g_{\omega\pi B} = 2g_{P' K\bar{K}} = 2g_{\omega K\bar{K}} = 2g_{P' Q\bar{K}} = 2g_{\omega Q\bar{K}} . \end{cases}$$

With these relations, only 2 residues in πN scattering are to be determined g^1 and g^0

$$g^1 = g_{\rho N\bar{N}} \; g_{\rho\pi\pi} \qquad\qquad g^0 = g_{P' N\bar{N}} \; g_{P'\pi\pi} .$$

Since there are two s-channel helicity amplitudes M_+ and M_- (see appendix 1) the number of independent parameters is actually 4.

The Pomeranchuck singularity is taken as a simple pole with a linear trajectory

$$P = i\sigma_\infty \; \nu \; e^{at} \; \nu^{\alpha_P(t)-1} \; e^{-\frac{i\pi}{2}(\alpha_P(t)-1)} . \tag{15}$$

b) Regge expansion in πN scattering

The amplitudes for πN scattering are defined in terms of t channel isospin, M^0 and M^1 :

$$\begin{aligned} \pi^- p &\to \pi^- p \quad : \quad M^0 + M^1 \\ \pi^+ p &\to \pi^+ p \quad : \quad M^0 - M^1 \\ \pi^- p &\to \pi^0 n \quad : \quad -\sqrt{2}\, M^1 \\ \pi^- p &\to \eta n \quad : \quad \frac{1}{\sqrt{3}}\, M'^1 . \end{aligned}$$

The amplitude M'^1 is deduced from M^1 by changing the signature of all contributions.

Having parametrized the input Regge poles, one obtains the cuts corresponding to the diagrams of figure 2.3 by convolution. The different contributions of the expansion (1) are listed in table II. Introducing some

arbitrary weights in order to take into account the unknown strength of Regge cuts, we parametrize our amplitudes as follows:

$$M^{\circ} = P+P' * (1+2aP) + b \{(P'*P' - \omega*\omega) + \tau(A_2*A_2 - \rho*\rho)\} * (1+2aP) +$$

$$+ d^{\circ} \{(P'*P*P' + \omega*P*\omega) + \tau(A_2*P*A_2 + \rho*P*\rho)\}$$

$$M^{1} = \rho*(1+2aP) + 2b(\rho*P' - \omega*A_2) * (1+2aP) + 2d^{1}(\rho*P*P' + \omega*P*A_2) \qquad (16)$$

$$M'^{1} = A_2*(1+2aP) + 2b(A_2*P' - \omega*\rho) * (1+2aP) + 2d^{1}(A_2*P*P' + \omega*P*\rho) .$$

In expression (16), one did not take into account helicity. The complete expressions for helicity amplitudes are given in the appendix 2. Let us discuss now the physical meaning of the parameters we choose.

i) absorptive cuts

In agreement with the most commonly used reggeized absorption models, we use only two parameters : the strength of absorption in non-flip (a_+) and flip amplitudes ($a_- = \mu a_+$) .

It is the same for different isospin amplitudes and different absorptive cuts (R*P and R*R*P). The weak absorption model corresponds to $a_+ = 1$. This choice is made in our fit, in agreement with our theoretical arguments. (see part 2-2).

ii) R*R cuts

In equation (16), we have assumed that all different (R*R) duality diagrams have the same weight : we obtain the same contribution in flip and non flip helicity amplitudes, a factor 2 due to the symmetrization $(R_1*R_2+R_2*R_1)$ in M^1 and finally a determination of τ ($\tau = 11/3$: see appendix 3). (R*R) cuts cancel in the amplitude M_+^1 , but not in M_-^1 : this important point for the determination of the parameter b in the fit is shown as follows. From table II, we get :

$$(R*R) \propto r^{\circ}*\varphi^{1} - \varphi^{\circ}*r^{1} = g^{\circ}g^{1}\left[(\sqrt{-t})^{n}\, 1*e^{-i\pi\alpha} - 1*e^{-i\pi\alpha}\,(\sqrt{-t})^{n}\right] ,$$

where we used an abbreviate notation : $r^i = g^i\, 1$, $\varphi^i = g^i\, e^{-i\pi\alpha}$. In the non

flip amplitude ($n = 0$) , the result is zero. In the flip amplitude ($n = 1$), there is only a partial cancellation, as can be shown from the convolution formula of two different exponentials in the t variable :

$$(\sqrt{-t})e^{\beta t} * e^{\beta' t} - e^{\beta t} * e^{\beta' t}(\sqrt{-t}) = \frac{\sqrt{-t}}{16\pi s}\, e^{\beta\beta' t/\beta+\beta'} \times \frac{\beta-\beta'}{(\beta+\beta')^2} \neq 0$$

with here $\beta = \text{cste} + \alpha' \log s$, $\beta' = \text{cste} + \alpha' (\log s - i\pi)$.

In practical calculations we take $|b| = 1/2$, which corresponds to the strength ± 1 for each duality-diagram contribution. The sign, undetermined from any of our theoretical arguments is determined from the phenomenological analysis.

iii) R*P*R cuts

Our parameters are d^1 and d^0 , the strength of R*P*R cuts in isospin 1 and isospin 0 amplitudes. In the flip amplitudes, no new parameter is introduced, the same ratio as for absorption corrections having been choosen for flip over non flip strengths of the cut.

$$\frac{d^1_-}{d^1_+} = \frac{d^0_-}{d^0_+} = \mu \left(= \frac{a_-}{a_+} \right).$$

Our phenomenological analysis shows that d^0 and d^1 cannot be taken equal, in particular as concerns the cross-over in πN and KN reactions. It appears that the same amount of $\omega * P' * P'$ and $\rho * P * P'$ is obtained only with $d^0 = 2d^1$ (see appendix 2). We considered it as an empirical rule, justified by a remark made at the end of the section 2.2 : diagonal R*P*R diagrams- such as the diagram $(\omega+P')*P*(\omega+P')$ - can be strongly enhanced with respect to non diagonal ones - such as $(\rho+A_2)*P*(P'+\omega)$-.We are then left with only one new parameter $d = d^1 = d^0/2$.

d) Regge expansion in KN scattering

The main difference is the absence of R*R and R*R*P cuts (see section 2.1). In table II are shown the contributions of the other Regge cuts for the following amplitudes describing all the reactions we consider :

$$
\begin{aligned}
K^+p &\to K^+p \qquad M = M^0_{ex} + M^1_{ex} \\
K^-p &\to K^-p \qquad M = M^0_{nex} + M^1_{nex} \\
K^+n &\to K^0p \qquad M = \quad 2M^1_{ex} \\
K^-p &\to \bar{K}^0n \qquad M = \quad 2M^1_{nex}
\end{aligned}
$$

the amplitudes of exotic (M_{ex}) and non exotic channels (M_{nex}) are related by s-u crossing. SU_3 and exchange degeneracy relations (14) allow to express directly the amplitudes of KN scattering in terms of the parameters of πN scattering, if we take the same parameters d^0 and d^1. We get

$$
\begin{aligned}
M^0_{ex} &= P + \tfrac{1}{2}(P'+\omega)*(1+2aP) + \frac{d^0}{2}\left[(P'+\omega)*P*(P'+\omega) + \tau'(A_2+\rho)*P*(A_2+\rho)\right] \\
M^0_{nex} &= P + \tfrac{1}{2}(P'-\omega)*(1+2aP) + \frac{d^0}{2}\left[(P'-\omega)*P*(P'-\omega) + \tau'(A_2-\rho)*P*(A_2-\rho)\right] \\
M^1_{ex} &= \tfrac{1}{2}(A_2+\rho)*(1+2aP) + d^1(P'+\omega)*P*(A_2+\rho) \\
M^1_{nex} &= \tfrac{1}{2}(A_2-\rho)*(1+2aP) + d^1(P'-\omega)*P*(A_2-\rho).
\end{aligned}
\qquad (17)
$$

In the expressions (17), all Reggeons (except P) have the same couplings as in $\pi N \to \pi N$ reactions. The only change with relations (16) is τ', equal to 5 (see appendix 3).

We note that the strength of R*P*R cuts, remains the same (with respect to Regge poles): as an example compare M^1_{ex} (formula 17) to the combination $(M^1+M'^1)/2$ (formula 16). This is not an obvious result, since it comes from a cancellation of two factors : The SU_3 factor 1/2 which relates $\pi\pi$ couplings to $K\bar{K}$ couplings (formula 14) and a factor 2 due to the contribution of K and Q intermediate states in all RPR cuts.

3.2. Results

The parameters of the model are given in table III. We have 12 free parameters, including those of Regge trajectories. The starred parameters in the table are fixed "a-priori" from theoretical considerations.

a) Calculations and results

The 12 parameters are determined, using a set of data on pion-nucleon reactions. The fit leads to excellent results for other pion-nucleon observables and all kaon-nucleon data in the usual range of validity for reggeized models ($5 \leq P_{lab} \leq 30$ GeV/c , $0 \leq -t \leq 1(GeV/c)^2$).

i) The first step was to fix all parameters of Regge poles (table III.a) with a fit to the total cross sections ($\sigma_{\pi^\pm}$, fig. 3 .1) and the differential cross-section ($d\sigma_{\pi^\pm}$, $d\sigma_x$, fig. 3 .3,4) in the energy range $5 \leq P_{lab} \leq 20$ GeV/c . As it is well-known, the ratio $g^1_- / g^1_+ = -4$, is taken from the ratio of the electromagnetic form factors of the nucleon.

ii) The next step consisted in fixing the Regge cut parameters, μ and d, (table III.b), using all our set of data, including elastic polarizations ($P^\pm$, fig. 3 .5) and imposing the position of the cross-over point, as determined by a recent accurate measurement of $d\sigma_{\pi^\pm}$ at low energy[17]. Note that there is another "free" parameter namely the sign of dual (R*R) cuts. The sign is determined from elastic polarizations which are sensitive to such a small contribution near the symmetrical zero ($-t \simeq .6\,(GeV/c)^2$) .

iii) The curves for the other observables in pion-nucleon scattering are predictions of the model : the rotation parameters (R,A,fig. 3 .6), the charge exchange polarization parameter (P^x, fig. 3 .8). It is well-known that absorption corrections give a non-zero P^x , but negative after the transfer momentum value of the cross-over[18]. In our model, we obtain a polarization staying positive at this point, in good agreement with data(*)[19]. More generally, the results of our fit in pion-nucleon scattering are excellent, showing that we reach a good description of the amplitudes themselves. This is well confirmed by the comparison of our amplitudes with the results of the amplitude analysis at $P_{lab} = 6$GeV/c (fig. 3 .9,10).[20]

(*) Our agreement is better with the Cern data than with the Argonne data. It should be remarked that this is not a test for choosing the best data or the best model since the same qualitative features are present in both sets of data : a non negative polarization, with a variation of the peak with incident energy. In particular a smaller value of d (not choosen by the minimization) would agree with the Argonne data.

iv) The reaction $\pi^- p \to \eta n$ is predicted, using SU_3 invariance and the experimental value of the $\Gamma(\eta \to 2\gamma)/\Gamma(\eta)$ branching ratio[21]: The differential cross section predictions ($d\sigma^\eta$ fig. 3 .7) and the polarization (P^η, fig. 3 .8) are in satisfactory agreement with data.

v) In the KN reactions we have to fix the 2 parameters of the Pomeranchuck singularity (table III.a). A simple geometrical relation reduces this number to one parameter

$$\frac{(\sigma_\infty)_{\pi N}}{a_{\pi N}} = \frac{(\sigma_\infty)_{KN}}{a_{KN}}$$

Then, with only one parameter, assuming SU_3 invariance and exchange degeneracy, all is predicted in the four studied reactions : the total cross-sections (figure 3 .11), the elastic (fig.3.12) and charge exchange differential cross-sections (fig. 3 .13) and polarizations (fig. 3 .14a,b). The good results we obtain are confirmed by the analysis of $\bar{K}N$ amplitudes[22] (fig. 3 .16). Note in particular the good results for the polarization in $K^- p \to K^o n$ (fig. 3 .15). This is a pure prediction since the (preliminary) data at $P_{lab} = 8$ GeV/c have just now become available[23] after we obtained our results.

b) Sensitivity of the fit and the role of secondary cuts

The sensitivity of the fit to the Regge-poles parameters is strong, due to their small number. Only slight changes of $(\sigma_\infty)_{\pi N}$, $(\sigma_\infty)_{KN}$ g_+^o , g_+^1 , and α_o are allowed the parameters α',α'_p ,$(a)_{\pi N}$ and ν_o can vary a little more, since they are important only in the t-dependence of the cross-sections. The parameter ν_o , in particular, is fixed by the differential cross section $d\sigma_x$ at larger transfer. It corrects the behaviour of all Regge residues together. The less important parameter is g^o , limited to small values . In particular, no clear choice for its sign seems to be indicated by data (we differ on that point from amplitude analysis.See figure 3 .9).

The Regge cut parameters (table III.b) are less precisely fixed by the data, since the gross features of the data are well reproduced by the Regge poles. Nevertheless their order of magnitude and their sign are well fixed by the role they play as far as the detailed structures of the data are concerned, making it possible to go from a qualitative to a quantitative understanding of the experimental results.

i) The absorption strength

The parameter a_+ was fixed at one in our fit, as previously indicated. But relaxing this constraint does not change the final result : a minimization gave $a_+ = .99$. Let us insist on this point, which is an important feature of the model. It is known that the physical amplitudes with negative signature have peripheral imaginary parts[24]. This is also true for our amplitudes as shown in the impact parameter representation (fig. 3.17). In our "weak" absorption model, R*P cuts alone are not strong enough to obtain peripherality, indicated for instance by the position of the zero in Im M^1_+ ($\simeq F^1_{//}$, see fig. 3.10). The R*P*R cut gives in the imaginary parts the amount of absorption that is lacking. This is opposite to what is called the strong absorption mechanism[9] for which all the peripheral character is due to R*P cuts.

The value of μ ($\mu = a_-/a_+$) which is mainly sensitive to σ_x and $P^{\pm}$ at larger momentum transfer ($-t \geq -.5 (GeV/c)^2$ demonstrates that absorption corrections have to be weaker in flip amplitudes.

We now turn to the role played by the different cut contributions.

ii) The R*P*R cut

The R*P*R cut plays a key role in the non flip amplitudes (the R * P * R cuts are also depressed by the factor μ , as well as absorption corrections).

- In negative signature amplitudes (ρ,ω exchange) it increases the absorption in the imaginary part, as previously mentioned and decreases the absorption in the real part, introducing a noticeable central contribution (fig. 3.17). This can be seen in another way through the zero structure of the amplitude M^1_+ , where, contrary to the usual absorption and other peripheral models, $Re M^1_+$ vanishes only after $Im\, M^1_+$[5] (compare $F^1_\perp$ to $F^1_{//}$ in figures 3.10). This leads to the good sign of the polarization P^x (Fig. 3.8) and the correct cross-over at low energy in $\pi^{\pm} p \to \pi^{\pm} p$ and $K^{\pm} p \to K^{\pm} p$ as shown by the prediction for the quantity Δ at $P_{lab} = 5$ GeV/c (fig. 3.18)

$$\Delta = \frac{[(d\sigma)_- - (d\sigma)_+]}{\sqrt{8[(d\sigma)_- + (d\sigma)_+]}}$$

- In positive signature amplitudes, we have the opposite effect : the real part is more absorbed and the imaginary part less. This gives a striking violation of exchange degeneracy. In particular one may note the double zero in $\mathrm{Im}P'$, in contrast with the single zero in $\mathrm{Im}\rho$ (fig. 3.19) and obtained without there being any zero in the Reggeon residue[25]. In the same way, a violation of exchange degeneracy appears in the low energy extrapolation for the KN and $\bar{K}N$ charge exchange differential cross-sections (fig. 3.20). In fact, in qualitative agreement with the data[26], the exotic cross-section ($K^+n \to K^\circ p$) is larger than the non exotic one ($K^-p \to \bar{K}^\circ n$). Note also the similarity with the hypercharge exchange cross-sections[27], for which the absorption model gives the opposite prediction.

Another effect of R*P*R cuts, is to give a small break in the elastic differential cross-sections, since it corresponds to a small, but slowly decreasing function of t . This effect near $-t = .8(\mathrm{GeV/c})^2$ is too small to agree with the data but its energy dependence is in agreement with them, showing that a better treatment of large transfer in the parametrization of the Regge poles for instance, can easily improve the fit.

iii) The R*R cuts

As previously mentioned R*R cuts contribute in the M_-^1 amplitude but not in M_+^1 . In fact, they play a role in P^x , avoiding a zero at $-t = .6(\mathrm{GeV/c})^2$, which would exist if the Regge poles only or with weakly coupled R*P and R*P*R cuts were present. This fixes their sign but their strength is not accurately determined.

c) Predictions and limitations of the model

Our set of results shows a good quantitative agreement with the data over the range in energy and transfer which is usually considered in reggeology. In particular, when compared with the usual absorption models, our result is that the phases and the zeros of the amplitudes, that is the t-dependence, are essentially well reproduced in the domain $-t \leq 1(\mathrm{GeV/c})^2$. The question of s-dependence is more controversal and is to be discussed now in more detail.

At very high energy, ($p_{lab} \geq$ 20-30 GeV/c) one expects that the R and R*P contributions should dominate the secondary contributions R*R , R*R*P and R*P*R , due to the difference in intercept (at least at small t). This means that all features of a pure absorption model will appear at high enough energy. In particular the cross-over point is expected to go to larger value of t (fig. 3 .21) and the charge exchange polarization to get a zero at small

transfer (fig. 3 .22). We believe that the second effect is a better test of the validity of our model than the first one, due to the fact that it depends on the real part of the amplitude, which is less affected than the imaginary part by possible changes in the Pomeranchuck amplitude or in the strength of absorptive cuts. We then propose to experimentalists measurements of P^x at an impulsion in the laboratory higher than 20 GeV/c.

At small energy (P_{lab} < 6 GeV/c) one expects the effects of the secondary cuts to increase. Almost qualitative agreement is obtained with πN differential cross sections at P_{lab} =2.5 GeV/c (fig. 3.3 and 4). In KN scattering, one obtains a qualitative agreement with the low energy data on charge exchange differential cross-sections, as already mentioned (fig. 3 .20).

Nevertheless, one problem arises about the question of s-dependence : at low energy and near the forward direction, the R*P*R cut seems to give a too large contribution to the K^+p total cross - section under P_{lab} = 8 GeV/c (fig. 3.11). This fact is related to the detailed F.E.S.R. analysis of Worden[28], which shows that the imaginary part of the non flip amplitude is too much absorbed at low energy, due to the stronger effect of the R*P*R cut. In fact all models which propose to modify the absorption model in order to improve it have the same difficulty. That is the reason why we think that the problem is not due to the R*P*R cut but to other singularities contributing to the imaginary part, probably the Pomeranchuck singularity. Other results (very high energy differential and total pp cross-sections, for instance) show that the Pomeranchuck singularity is more complicated than a simple Regge pole. In fact a dual Pomeranchuck singularity corresponding to Regge cuts in the s channel, would have to be smaller than the actual pole, at low energy, since the cuts give an increasing contribution to the background at small s , instead of a constant one. An explicit example of such a Pomeranchuck singularity is given in R.Lacaze's thesis[29]. Then the R*P*R cut is not obliged to be small, and could give the precise contribution which has to be added at low energy to an increasing but small Pomeranchuck contribution.

4 - CONCLUSION AND OUTLOOK

Though it is based on a small modification of the reggeized absorption models since it consists in taking into account secondary cuts, our model leads to some interesting conclusions as far as the theory and the phenomenology related

to duality problems is concerned.

a) Phenomenological conclusions and tests of the model

First, the smallness of the modification itself makes one believe in the basic validity of the absorption idea, as applied to Regge poles. R*R and R*P*R cuts are naturally obtained as corrections to the first dominant Regge singularities (P,R,R*P) which make up the reggeized absorption model. Secondly, duality makes it possible to find selection rules which help in reducing the number of free parameters. Let us recall here the main features of R*P*R cuts :

i) they contribute to exotic channels

ii) they strongly contribute in negative signature amplitudes, increasing the absorption in the imaginary part, and decreasing it in the real part of the amplitude

iii) they give a violation of exchange degeneracy, since in the positive signature amplitude they reduce the absorption in the imaginary part contrary to what happens in the negative signature amplitude. R*R cuts, for which duality selection rules forbid those same effects, nevertheless play a role in some amplitudes in particular in that with s-channel helicity flip and negative signature. All the secondary cut contributions modify the amplitudes of the absorption model in such a way as to give a quantitative description of all observables in πN and KN scattering, in particular the charge exchange and elastic polarizations and the cross-over effect.

The question arises of finding a test of the model, especially of the secondary Regge cuts we introduce. One test is given by the energy variation of the model amplitudes : at high energy, a typical property of the model is to tend to the "weak" absorption model itself, with a charge exchange polarization changing sign at some high energy (see fig. 3.22). A second test can be provided by studying other reactions where phenomenological problems appear. In particular, the violation of exchange degeneracy in hypercharge exchange reactions[27] ($K^-p \to \pi^-\Sigma^+$, $\pi^+p \to K^+\Sigma^+$) gives identical effects as in low energy KN and $\bar{K}N$ reactions ($K^-p \to \bar{K}^0 n$, $K^+n \to K^0 p$) . The comparison with data in these reactions, by a straightforward application of our model, will then be very useful. The extension of our model to meson-meson and Baryon-Baryon reactions (with different duality diagrams and selection rules) is also under study. A third test concerns more particularly R*R cuts, which are less sensitive to the fit in the reactions we have studied. R*R cuts intervene in the description

of forbidden peaks or t-channel exotic reactions. Our selection rules (see table I) forbid R*P*R cuts in such reactions, but not R*R cuts. Therefore we consider it very important to study these reactions in connection with the particular selection rules of R*R cuts.

b) Duality, unitarity and a theoretical outlook

The Regge cuts play in our model a very important role within the framework of duality. The R*P*R cut gives a natural mechanism for breaking exchange degeneracy, in contrast with a model involving Regge poles only as in the first versions of duality. Exchange degeneracy appears as an exact symmetry property of Regge poles, with a "weak" violation due to absorption corrections (if ReP $\neq$ 0) and a "maximal" violation by R*P*R cuts (the R*P*R contribution in an exotic s-channel is nearly imaginary).

Moreover incorporating R*P*R and R*R cuts into phenomenological duality relation between resonances and t-channel singularities (now Regge poles and cuts) leads to a new aspect of this relation which can be called "duality for real parts". Indeed it is possible to associate the central contribution in the real parts of the amplitudes - characteristic of R*P*R and R*R cuts (fig. 3.17) - with similar features of low energy amplitudes due to resonances. We know that the main resonances are peripheral, since their physical spin J and masses M obey the empirical relation

$$b = \frac{J + \frac{1}{2}}{M} \simeq 1 \text{ fermi} ,$$

where b is the impact parameter. In fact the imaginary parts of the peripheral resonances are strong near the corresponding energy ($\sqrt{s} = M$), but the real parts vanish at this point and have extrema at lower and larger $\sqrt{s}$ (cf. the Breit-Wigner formula). In terms of the impact parameter, this gives respectively larger and lower values than 1 fermi, that is non peripheral real parts. It is possible to show by comparison of the amplitude analysis in the resonance domain ($P_{lab} = 2.5$ GeV/c) and the Regge domain ($P_{lab} = 6$ GeV/c), that the two effects are related.[20]

The use of s-channel unitarity (cf. section 3.2) provides theoretical arguments for R*P and R*P*R cuts. It connects them and gives satisfactory results for their signs and strengths except for the flip amplitudes, for which an explanation of why they are practically not absorbed is still lacking (cf. however Lovelace[30]). The modification proposed to the multiperipheral graph is one of the first corrections which appear in a dual perturbative expansion

of many particle production amplitudes. Such an approach to many body production processes, together with the same for two-body scattering, seems to be an interesting open field for phenomenology.

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APPENDIX 1

KINEMATICS AND NOTATIONS

KINEMATIC VARIABLES

We consider two-body reactions, which are described by amplitudes depending on the variables s,t,u : (figure A-1)

$$s = (p_1 + p_2)^2 \quad , \quad t = (p_1 - p_3)^2 \quad , \quad u = (p_1 - p_4)^2$$

$$p_1^2 = p_3^2 = m_1^2 \quad , \quad p_2^2 = p_4^2 = M^2 \quad , \quad s + t + u = 2M^2 + 2m^2$$

One also uses the variable $\nu = (s-u)/2$

We use kinematical variables defined in the center-of-mass frame - the momentum p, the scattering angle θ -, or in the laboratory frame - the energy and momentum of the incident particle E_{lab} , P_{lab} -. We have the relations :

$$p = \frac{1}{2\sqrt{s}} \left[\left(s - (M-m)^2\right)\left(s - (M+m)^2\right) \right]^{1/2} \quad , \quad \cos\theta = 1 + \frac{t}{2p^2}$$

$$E_{lab} = \frac{s - M^2 - m^2}{2M} \qquad P_{lab} = \left(E_{lab}^2 - m^2\right)^{1/2}$$

AMPLITUDES AND OBSERVABLES

The reactions we consider are of the type (spin 0+spin 1/2 → spin 0 + spin 1/2). The scattering amplitudes we use are s-channel helicity amplitudes. We have four amplitudes $M_{\lambda',\lambda}$, depending on the helicities of the nucleon $\lambda,\lambda' = \pm \frac{1}{2}$, among which only two are independent, M_+ and M_-

$$M_{1/2,1/2} = M_{-1/2,-1/2} = M_+ \qquad \text{no helicity flip}$$

$$M_{1/2,-1/2} = -M_{-1/2,1/2} = M_- \qquad \text{helicity flip}$$

The observables are obtained using the following formulae

total cross-section : $\sigma_{tot} = \frac{1}{2p\sqrt{s}} \, Im \, M_+ (s,0)$

differential cross-section : $\frac{d\sigma}{dt} = \frac{1}{64\pi s p^2} \left(|M_+(s,t)|^2 + |M_-(s,t)|^2 \right)$

polarization : $P = 2 \, Im \, M_+ M_-^* / (|M_+|^2 + |M_-|^2)$

PARTIAL WAVE EXPANSION AND CONVOLUTION

The partial-wave expansion is written as follows

$$M_{\lambda'\lambda} = \sum_J (J + 1/2) \, d^J_{\lambda'\lambda}(\theta) \, M^J_{\lambda'\lambda}$$

where $d^J_{\lambda'\lambda}$ are the rotation matrix elements.

We define the convolution operator * on the helicity amplitudes

$$\left(M^1 * M^2\right)^J_{\lambda'\lambda} = \frac{ip}{8\pi\sqrt{s}} \sum_{\mu=\pm 1/2} \left(M^1\right)^J_{\lambda'\mu} \cdot \left(M^2\right)^J_{\mu\lambda}$$

APPENDIX 2

HELICITY AMPLITUDES OF THE MODEL

In this appendix, we establish the list of all helicity amplitudes for πN and KN scattering, in their definitive form for practical use. We consider the Regge pole amplitudes with their couplings to πN reactions (formula 14) and the 5 Regge cut parameters a, b, d^0, d^1, μ ($\tau = 11/3$, $\tau' = 5$, cf. app. 3).

For helicity amplitudes, one must take into account the summation over intermediate helicity in the computation of Regge cuts. The Pomeranchuck pole giving no helicity flip, only Regge poles are considered in the following formula

$$(A * B)_+ = (A * B)_{+\frac{1}{2}\,+\frac{1}{2}} = A_{+\frac{1}{2}\,+\frac{1}{2}} * B_{+\frac{1}{2}\,+\frac{1}{2}} + A_{+\frac{1}{2}\,-\frac{1}{2}} * B_{-\frac{1}{2}\,+\frac{1}{2}} = A_+ * B_+ - A_- * B_-$$

$$(A * B)_- = (A * B)_{+\frac{1}{2}\,-\frac{1}{2}} = A_{+\frac{1}{2}\,+\frac{1}{2}} * B_{+\frac{1}{2}\,-\frac{1}{2}} + A_{+\frac{1}{2}\,-\frac{1}{2}} * B_{-\frac{1}{2}\,-\frac{1}{2}} = A_+ * B_- + A_- * B_+$$

The πN amplitudes are written as follows :

$$M^0_+ = P + P'_+ * (1 + 2\alpha P) + b\{(P'_+ * P'_+ - \omega_+ * \omega_+) + \tau(A_{2+} * A_{2+} - \rho_+ * \rho_+)$$
$$- (P'_- * P'_- - \omega_- * \omega_-) - \tau(A_{2-} * A_{2-} - \rho_- * \rho_-)\}(1 + 2\alpha P)$$
$$+ d^0\{(P'_+ * P * P'_+ + \omega_+ * P * \omega_+) + \tau(A_{2+} * P * A_{2+} + \rho_+ * P * \rho_+)$$
$$- \mu^2(P'_- * P * P'_- + \omega_- * P * \omega_-) - \tau\mu^2(A_{2-} * P * A_{2-} + \rho_- * P * \rho_-)\}$$

$$M^0_- = P'_- * (1 + 2\alpha\mu P) + 2b\{(P'_+ * P'_- - \omega_+ * \omega_-) + \tau(A_{2+} * A_{2-} - \rho_+ * \rho_-)\} * (1 + 2a\mu P)$$
$$+ 2d^0\mu\{(P'_+ * P * P'_- + \omega_+ * P * \omega_-) + \tau(A_{2+} * P * A_{2-} + \rho_+ * P * \rho_-)\}$$

$$M^1_+ = \rho_+ * (1 + 2\alpha P) + 4d^1(\rho_+ * P * P'_+ - \mu^2\rho_- * P * P'_-)$$

$$M_-^1 = \rho_- * (1+2a\mu P) + 2b\{(\rho_+ * P_-' - \omega_+ * A_{2-}) + (\rho_- * P_+' - \omega_- * A_{2+})\} * (1+2a\mu P)$$
$$+ 2d^1\mu\{(\rho_+ * P * P_-' + \omega_+ * P * A_{2-}) + (\rho_- * P * P_+' + \omega_- * P * A_{2+})\}$$

$$M_+'^1 = A_{2+} * (1+2aP) + 2b\{(A_{2+} * P_+' - \omega_+ * \rho_+) - (A_{2-} * P_-' - \omega_- * \rho_-)\} * (1+2aP)$$
$$+ 2d^1\{(A_{2+} * P * P_+' + \omega_+ * P * \rho_+) - \mu^2(A_{2-} * P * P_-' + \omega_- * P * \rho_-)\}$$

$$M_-'^1 = A_{2-} * (1+2a\mu P) + 2b\{(A_{2+} * P_-' - \omega_+ * \rho_-) + (A_{2-} * P_+' - \omega_- * \rho_+)\} * (1+2a\mu P)$$
$$+ 2d^1\mu\{(A_{2+} * P * P_-' + \omega_+ * P * \rho_-) + (A_{2-} * P * P_+' + \omega_- * P * \rho_+)\}$$

The KN and $\bar{K}N$ amplitudes are expressed in terms of the same Regge pole amplitudes.

$$M_{ex+}^0 = P_{KN} + \frac{1}{2}(P_+' + \omega_+) * (1+2aP) + \frac{d^0}{2}\{(P_+' + \omega_+) * P * (P_+' + \omega_+)$$
$$+ \tau'(A_{2+} + \rho_+) * P * (A_{2+} + \rho_+) - \mu^2(P_-' + \omega_-) * P * (P_-' + \omega_-) - \mu^2\tau'(A_{2-} + \rho_-) * P * (A_{2-} + \rho_-)\}$$

$$M_{ex-}^0 = \frac{1}{2}(P_-' + \omega_-) * (1+2a\mu P) + \mu d^0\{(P_+' + \omega_+) * P * (P_-' + \omega_-) + \tau'(A_{2+} + \rho_+) * P * (A_{2-} + \rho_-)\}$$

$$M_{nex+}^0 = P_{KN} + \frac{1}{2}(P_+' - \omega_+) * (1+2aP) + \frac{d^0}{2}\{(P_+' - \omega_+) * P * (P_+' - \omega_+)$$
$$+ \tau'(A_{2+} - \rho_+) * P * (A_{2+} - \rho_+) - \mu^2(P_-' - \omega_-) * P * (P_-' - \omega_-) - \mu^2\tau'(A_{2-} - \rho_-) * P * (A_{2-} - \rho_-)\}$$

$$M_{nex-}^0 = \frac{1}{2}(P_-' - \omega_-) * (1+2a\mu P) + \mu d^0\{(P_+' - \omega_+) * P * (P_-' - \omega_-) + \tau'(A_{2+} - \rho_+) * P * (A_{2-} - \rho_-)\}$$

$$M_{ex+}^1 = \frac{1}{2}(A_{2+} + \rho_+) * (1+2aP) + d^1\{(P_+' + \omega_+) * P * (A_{2+} + \rho_+) - \mu^2(P_-' + \omega_-) * P * (A_{2-} + \rho_-)\}$$

$$M_{ex-}^1 = \frac{1}{2}(A_{2-} + \rho_-) * (1+2a\mu P) + 2\mu d^1\{(P_+' + \omega_+) * P * (A_{2-} + \rho_-) + (P_-' + \omega_-) * P * (A_{2+} + \rho_+)\}$$

$$M_{nex+}^1 = \frac{1}{2}(A_{2+} - \rho_+) * (1+2aP) + d^1\{(P_+' - \omega_+) * P * (A_{2+} - \rho_+) - \mu^2(P_-' - \omega_-) * P * (A_{2-} - \rho_-)\}$$

$$M_{nex-}^1 = \frac{1}{2}(A_{2-} - \rho_-) * (1+2a\mu P) + 2\mu d^1\{(P_+' - \omega_+) * P * (A_{2-} - \rho_-) + (P_-' - \omega_-) * P * (A_{2+} - \rho_+)\}$$

Let us check the empiral rule proposed in the text, relating d^0 and d^1. The contribution to negative signature in πN elastic scattering (non flip amplitude) is :

$$\rho_+ * (1 + 2\alpha P) + 4d^1 (\rho_+ * P * P'_+ - \mu^2 \rho_- * P * P'_-)$$

The same contribution in KN elastic scattering is mainly due to ω exchange ; we get from the negative signature contribution to M^0_{ex}:

$$\tfrac{1}{2} \omega_+ * (1 + 2\alpha P) + d^0 (\omega_+ * P * P'_+ + \tau' \rho_+ * P * A_{2+} - \mu^2 \omega_- * P * P'_- - \mu^2 \tau' \rho_- * P * A_{2-})$$

The main contributions of both amplitudes are

$$\rho_+ * (1 + 2\alpha P) + 4d^1 \rho_+ * P * P'_+ \quad \text{and} \quad \tfrac{1}{2}\left[\omega_+ * (1 + 2\alpha P) + 2d^0 \omega_+ * P * P'_+\right]$$

$d^0 = 2d^1$, is then the condition to have the same correction (up to an overall coupling) : this is taken as an empirical rule in our fit.

APPENDIX 3

CALCULATION OF THE FACTORS τ AND τ'

We compute the factors τ and τ' , which give the relative weight of the following contributions to isospin 0 amplitudes : double isospin 1 exchange and double isospin 0 exchange. τ and τ' are obtained by drawing all different R*R duality diagrams (see figure A.2 , the diagrams are drawn in a simplified but convenient form).

SU_2 and SU_3 Clebsh-Gordan coefficients make it possible to relate all the contributions of the diagrams . Applying the rule of the same weight for all diagrams, we get for the pion-nucleon case (fig.A.2 a,b,c,d)

$$\tau = \frac{(b) + (c) + (d)}{(a)} = \frac{1 + (\sqrt{2})^2 + (\sqrt{2})^2/(\sqrt{3})^2}{1} = \frac{11}{3}$$

In the Kaon nucleon case, we get (figure A. 2;a',b',c')

$$\tau' = \frac{(b') + (c')}{(a')} = \frac{1 + (\sqrt{2})^4}{1} = 5$$

(The same values τ and τ' were also used for R*R*P and R*P*R cuts in the amplitudes see appendix 2).

Note that in order to ensure correct signature properties of the Regge cuts, we are obliged to introduce crossed diagrams for some of them when there are neutral intermediate states (fig.A.3).

TABLE CAPTIONS

Table I : Classification and selection rules of duality diagrams.

Table II : Duality diagrams contributions in πN and KN amplitudes.

Table III: Parameters of the model.
The starred numbers correspond to the parameters we do not vary in the fit.

Duality Diagram	R	R*P	R*R	R*R*P	R*P*R
(s,t)	$e^{-i\pi\alpha}$	$e^{-i\pi\alpha} * P$			$e^{-i\pi\alpha} * P * e^{-i\pi\alpha}$
(u,t)	$\mathbb{1}$	$\mathbb{1} * P$			$\mathbb{1} * P * \mathbb{1}$
(s,u)			$e^{-i\pi\alpha} * \mathbb{1}$ $\mathbb{1} * e^{-i\pi\alpha}$	$e^{-i\pi\alpha} * \mathbb{1} * P$ $\mathbb{1} * e^{-i\pi\alpha} * P$	

TABLE I

πN scattering

	(P)	(R)	(R*P)	(R*R)	(R*R*P)	(R*P*R)
M^0	P	$r^0+\varphi^0$	$(r^0+\varphi^0)*P$	$r^0*\varphi^0$ $r^1*\varphi^1$	$r^0*\varphi^0*P$ $r^1*\varphi^1*P$	$r^0*P*r^0+\varphi^0*P*\varphi^0$ $r^1*P*r^1+\varphi^1*P*\varphi^1$
M^1		$r^1-\varphi^1$	$(r^1-\varphi^1)*P$	$r^0*\varphi^1-\varphi^0*r^1$	$r^0*\varphi^1*P-\varphi^0*r^1*P$	$r^1*P*r^0-\varphi^1*P*\varphi^0$
M'^1		$r^1+\varphi^1$	$(r^1+\varphi^1)*P$	$r^0*\varphi^1+\varphi^0*r^1$	$r^0*\varphi^1*P+\varphi^0*r^1*P$	$r^1*P*r^0+\varphi^1*P*\varphi^0$

KN scattering

M^0_{ex}	P	r^0	r^0*P		r^0*P*r^0 r^1*P*r^1
M^0_{nex}	P	φ^0	φ^0*P		$\varphi^0*P*\varphi^0$ $\varphi^1*P*\varphi^1$
M^1_{ex}		r^1	r^1*P		r^0*P*r^1
M^1_{nex}		φ^1	φ^1*P		$\varphi^0*P*\varphi^1$

TABLE II

a) Regge poles

	helicity	couplings	t-dependence	Regge trajectories
P	M_+	$(\sigma_\infty)_{\pi N} = 21.2$mb $(\sigma_\infty)_{KN} = 16.9$mb	$(a)_{\pi N} = 2.75$ $(a)^*_{KN} = 2.2$	$\alpha_p - 1 = .2t$
R	M_+	$g^0 = -19$ $g^1 = -4.5$	$S_o = .275$	$\alpha = .5 + 0.76t$
	M_-	$g^0 = -1$ $g^{1*} = 18$		

b) Regge cuts

	(R*P)	(R*R)	(R*R*P)	(R*P*R)
	a	b	c(=a×b)	d
M_+	1^*	0.5^*	1^*	- 3
M_-	0.5	0.5^*	0.5^*	-1.5^*

TABLE III

FIGURE CAPTIONS

Fig. 2.1 : Regge-Regge cut (a) with the sub-processes particle-Reggeon I and J of sub-energies s_i and u_i (b and c) and the corresponding dual representation (d)

Fig. 2.2 : Duality diagrams for Regge poles (a and b) and for the Pomeranchuck singularity (c and d)

Fig. 2.3 : Duality diagrams for Regge cuts : R*P(a) , R*R(b) , R*R*P(c) and R*P*R(d)

Fig. 2.4 : Multiperipheral graph and its Regge correction

Fig. 3 .1 : πN-total cross sections (data from ref.[31])

Fig. 3 .2 : Ratio $(\mathrm{Re}\, M_+ / \mathrm{Im}\, M_+)_{t=0}$ for $\pi^{\pm}p$ elastic scattering. The data are from ref.[32] (crosses for π^+p and circles for π^-p).

Fig. 3 .3 : $\pi^{\pm}p$ elastic differential cross section for different value of P_{lab} varying from 2.5.GeV/c to 50 GeV/c (data from ref.[33])

Fig. 3 .4 : Charge exchange differential cross section for P_{lab} varying from 2.5 GeV/c to 50 GeV/c . The data are from ref[34]

Fig. 3 .5 : Elastic polarization at P_{lab} = 6,10,14,40 GeV/c from ref.[35]

Fig. 3 .6 : R and A parameters for elastic and charge exchange πN scattering. The data are taken from ref.[36]

Fig. 3 .7 : Differential cross section for $\pi^-p \to \eta n$ from P_{lab} = 5.9 GeV/c up to P_{lab} = 18.2 GeV/c. The data are from ref.[21]

Fig. 3 .8 : Polarization for $\pi^-p \to \pi^0 n$ and $\pi^-p \to \eta n$ and the comparison with the CERN data[19] at P_{lab} = 4.9 GeV/c and 7.85 GeV/c

Fig. 3 .9 : Isospin zero amplitudes ($\sqrt{mb}/GeV/c$) versus -t in $(GeV/c)^2$. The flip amplitude is defined with respect to M°_{+} , $F_{//}$ being the parallel component and $F_{\perp}$ the orthogonal one to M°_{+} . Data are from ref.[20]

Fig. 3 .10 : Isospin 1 amplitudes with the same definition as in fig.9

Fig. 3 .11 : Total cross sections for $K^{\pm}$ p (data from ref.[37])

Fig. 3 .12 : $K^{\pm}$ p elastic differential cross sections (data from ref.[38])

Fig. 3 .13 : KN and $\bar{K}N$ charge exchange differential cross section. Data from ref [39]

Fig. 3..14 : $K^{\pm}$ p elastic polarization. Data from ref.[35a]

Fig. 3 .15 : $K^{-}p \rightarrow \bar{K}^{\circ}n$ polarization at $P_{lab} = 8$ GeV/c from ref.[23]

Fig. 3 .16 : M°_{+} and M^{1}_{-} in $\sqrt{\mu b}/GeV/c$ versus -t in $(GeV/c)^2$ for $\bar{K}N$ scattering the solid (dashed) line represents real (imaginary) part. Data from ref.[22]

Fig. 3 .17 : Hankel transform of isospin 1 amplitudes in arbitrary units versus the impact parameter in fermi units . The solid (dashed) line represents the real (imaginary) part.

Fig. 3 .18 : Difference between elastic scattering of particles and antiparticles from proton. We have quoted the quantity :

$$\Delta_x = \frac{\frac{d\sigma}{dt}(x^- p) - \frac{d\sigma}{dt}(x^+ p)}{\sqrt{8\left[\frac{d\sigma}{dt}(x^- p) + \frac{d\sigma}{dt}(x^+ p)\right]}} \quad \text{in } \sqrt{mb}/GeV/c$$

where x is π or K. Data from ref.[17]

Fig. 3 .19 : Comparison of imaginary parts of isospin 1 amplitude and non Pomeron isospin 0 amplitude (non flip parts) normalized to one in the forward direction

Fig. 3 .20 : Comparison of exotic (solid line) and non exotic (dashed line) KN charge exchange differential cross sections for P_{lab} varying from 1.2 GeV/c up to 100 GeV/c

Fig. 3 .21 : Variation of the cross-over point with energy

Fig. 3 .22 : Variation of the polarizations with energy ($\pi^- p \to \pi^o n$ and $\pi^- p \to \eta n$)

Fig. 3 .23 : Variation of the polarizations with energy ($K^- p \to \bar{K}^o n$ and $K^+ n \to K^o p$).

Fig.A.1 : Scattering 1+2 → 3+4

Fig.A.2 : Duality diagrams with Isospin 0 exchanged

Fig.A.3 : s-u crossed diagrams in the case of neutral intermediate states.

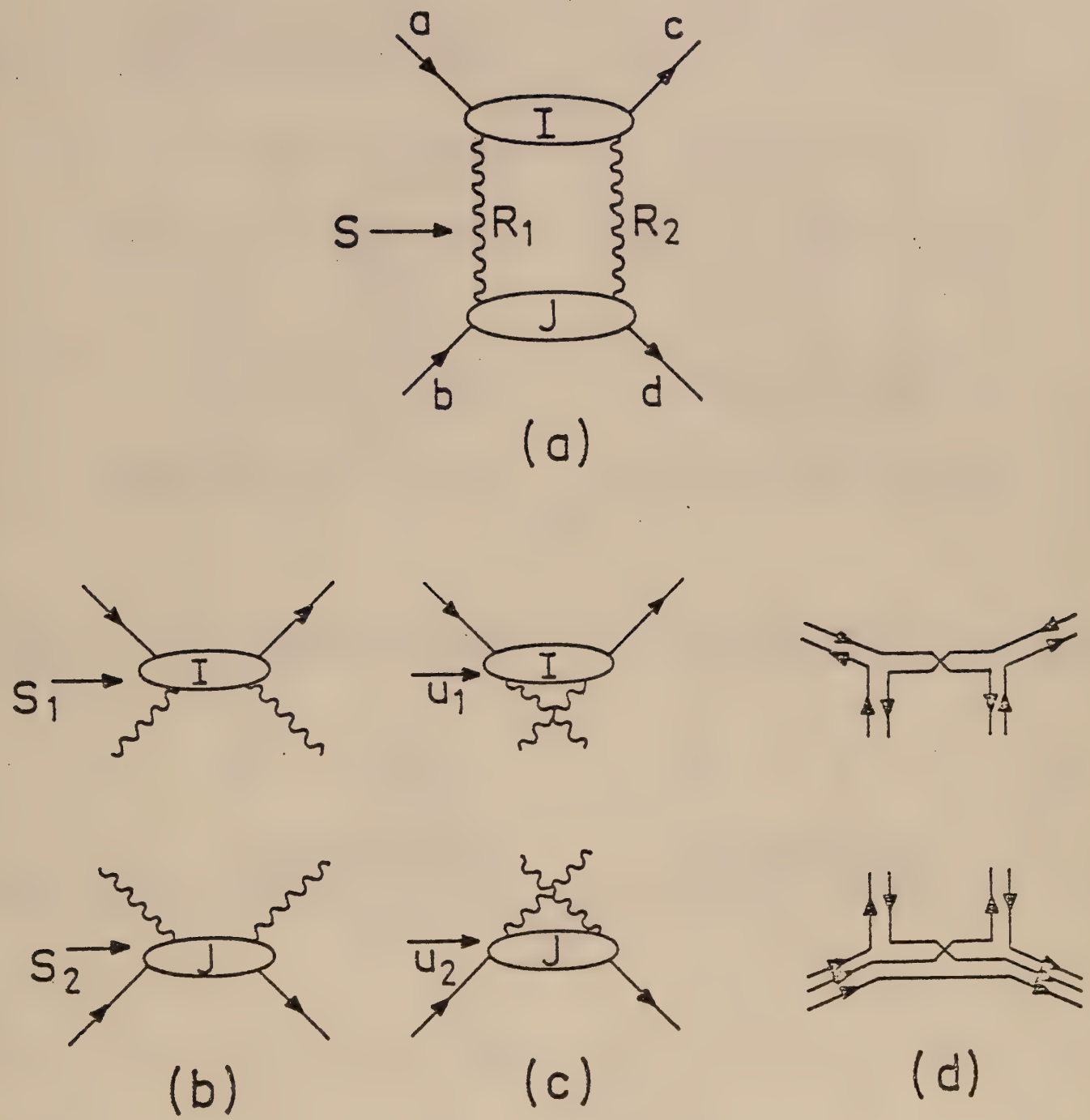


Fig.2.1

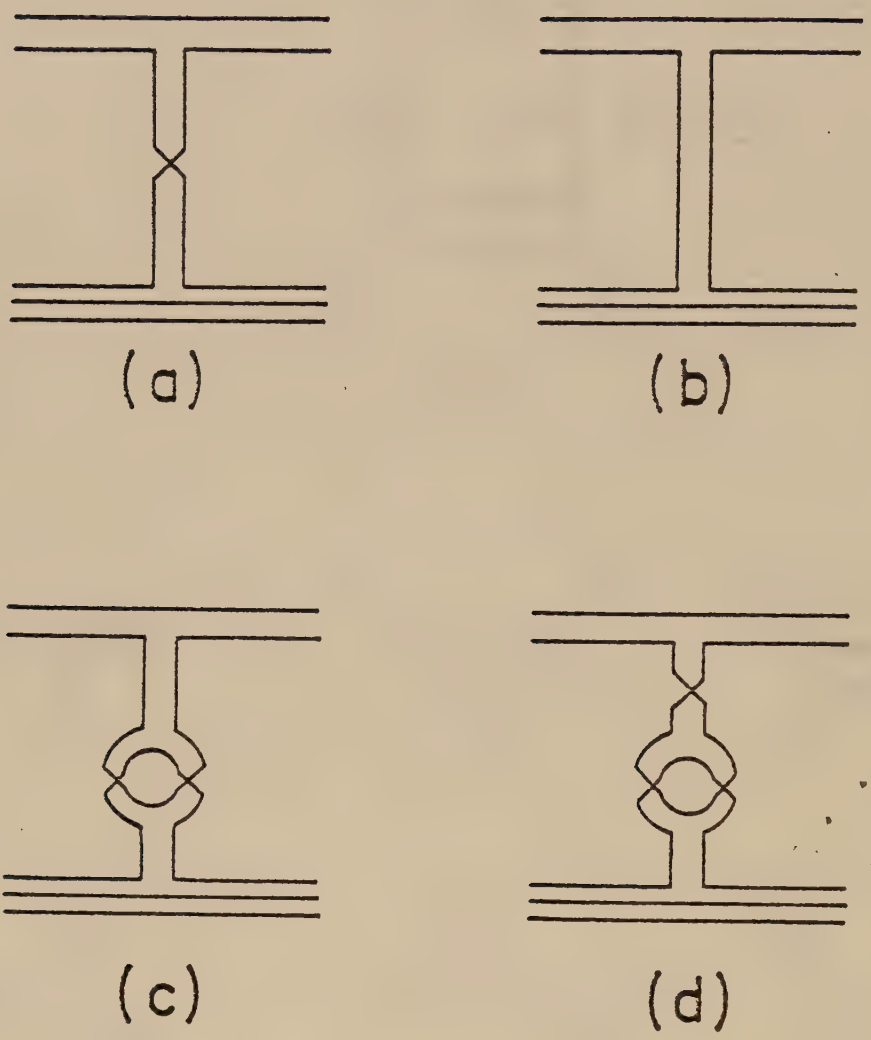
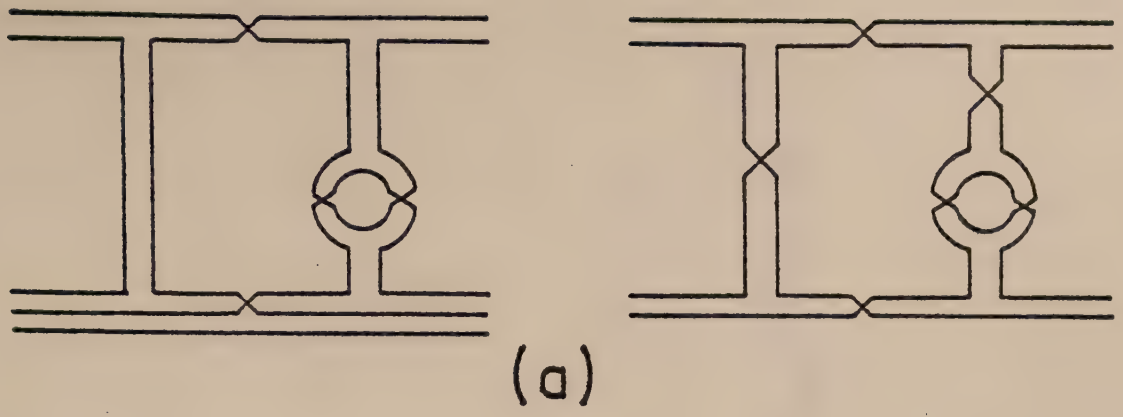
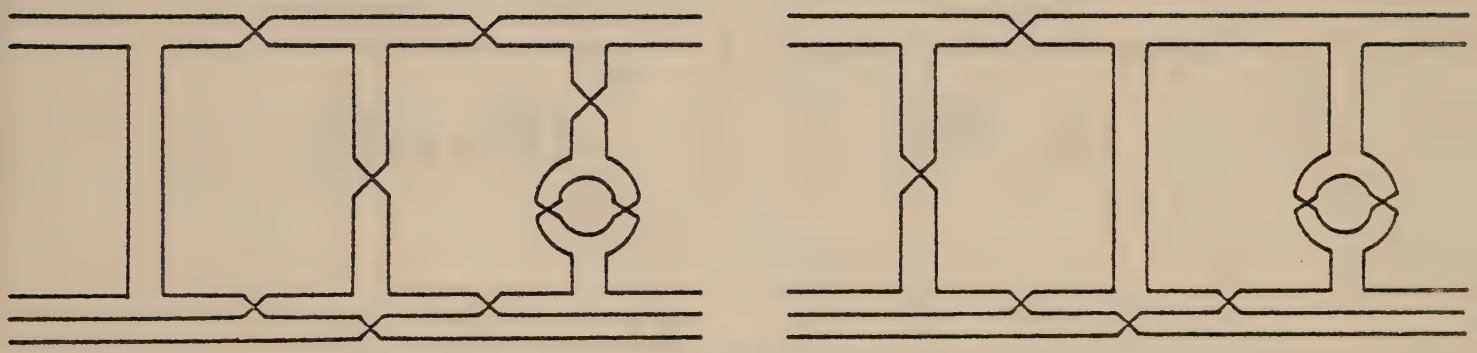


Fig.2.2

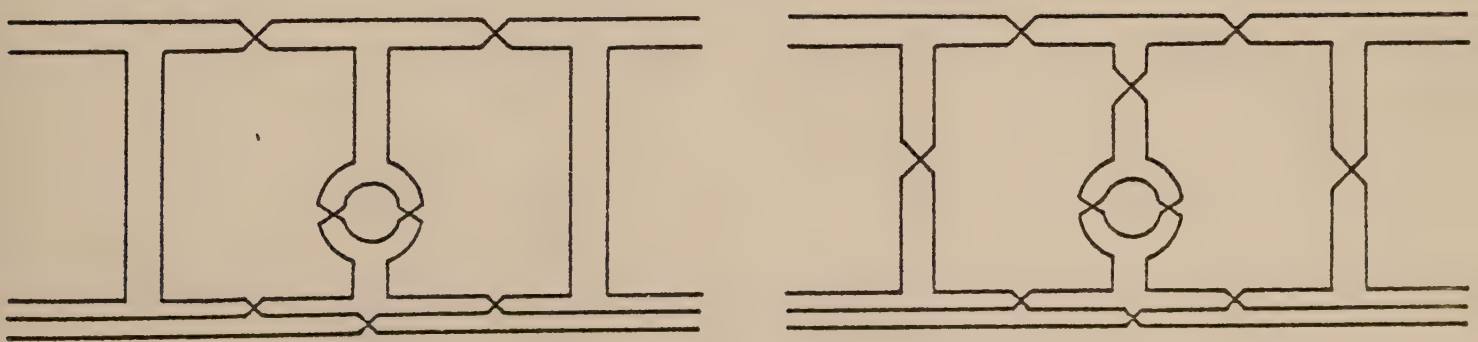


(a)

(b)

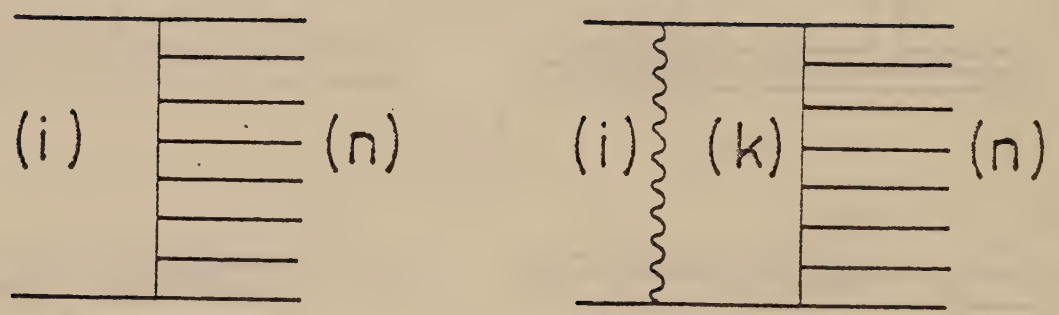


(c)



(d)

Fig.2.3



A_{ni} $(R * A)_{ni}$

Fig.2.4

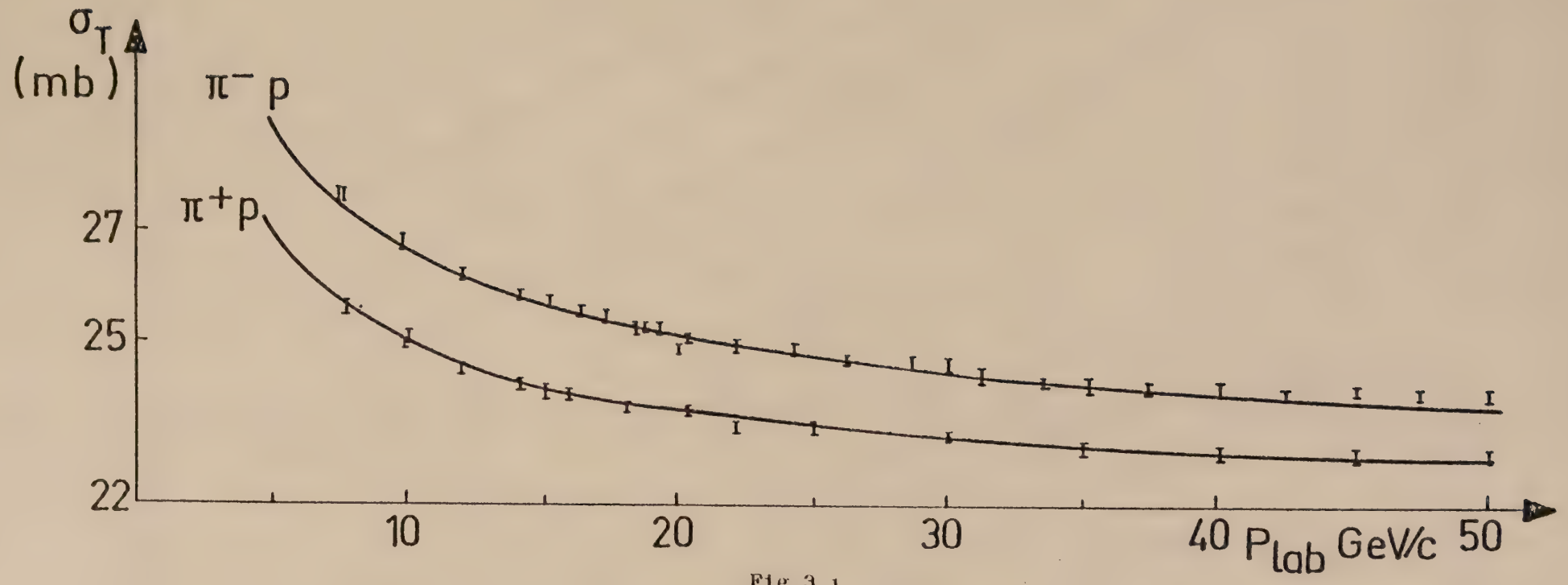
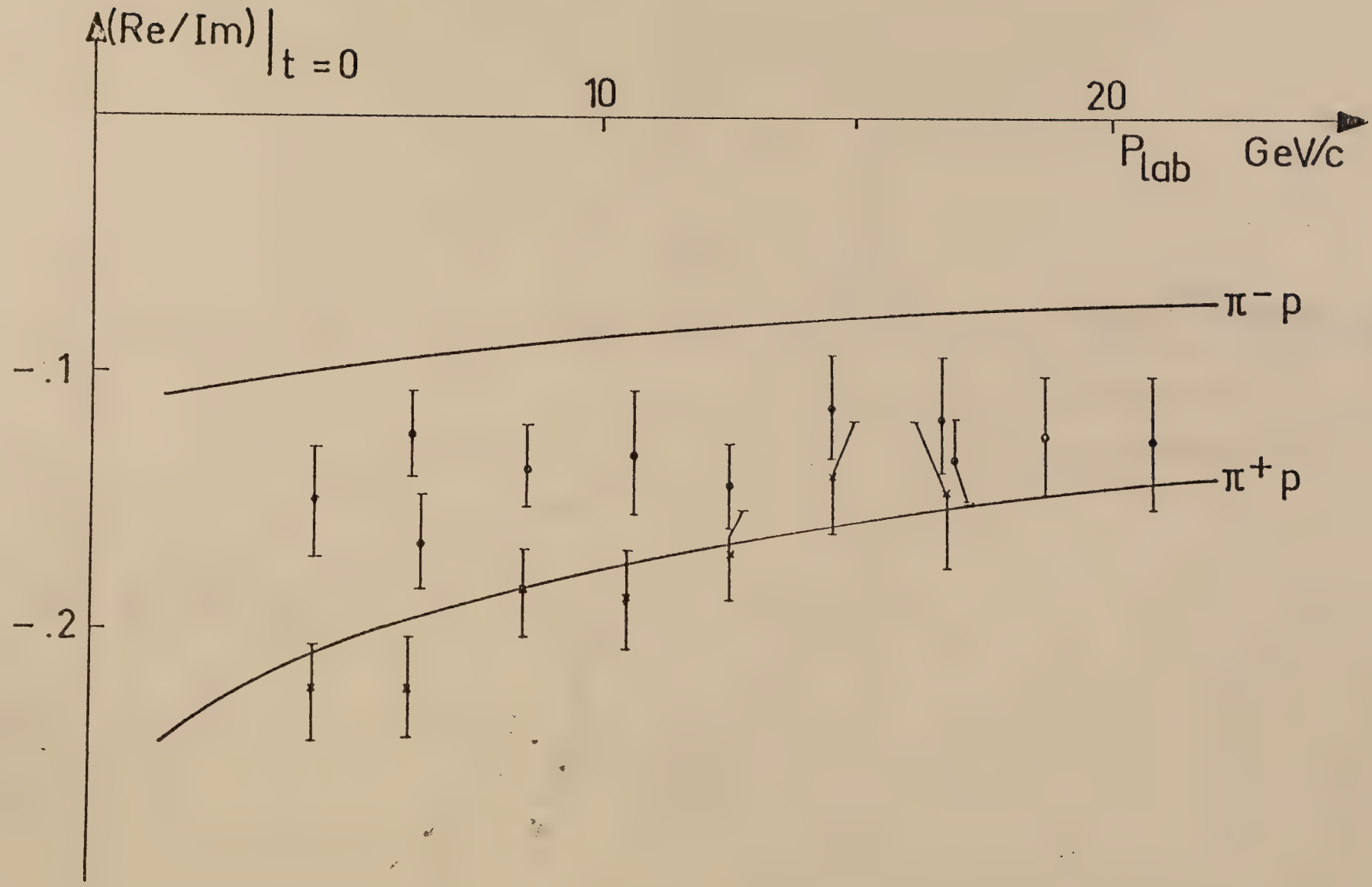


Fig.3.1



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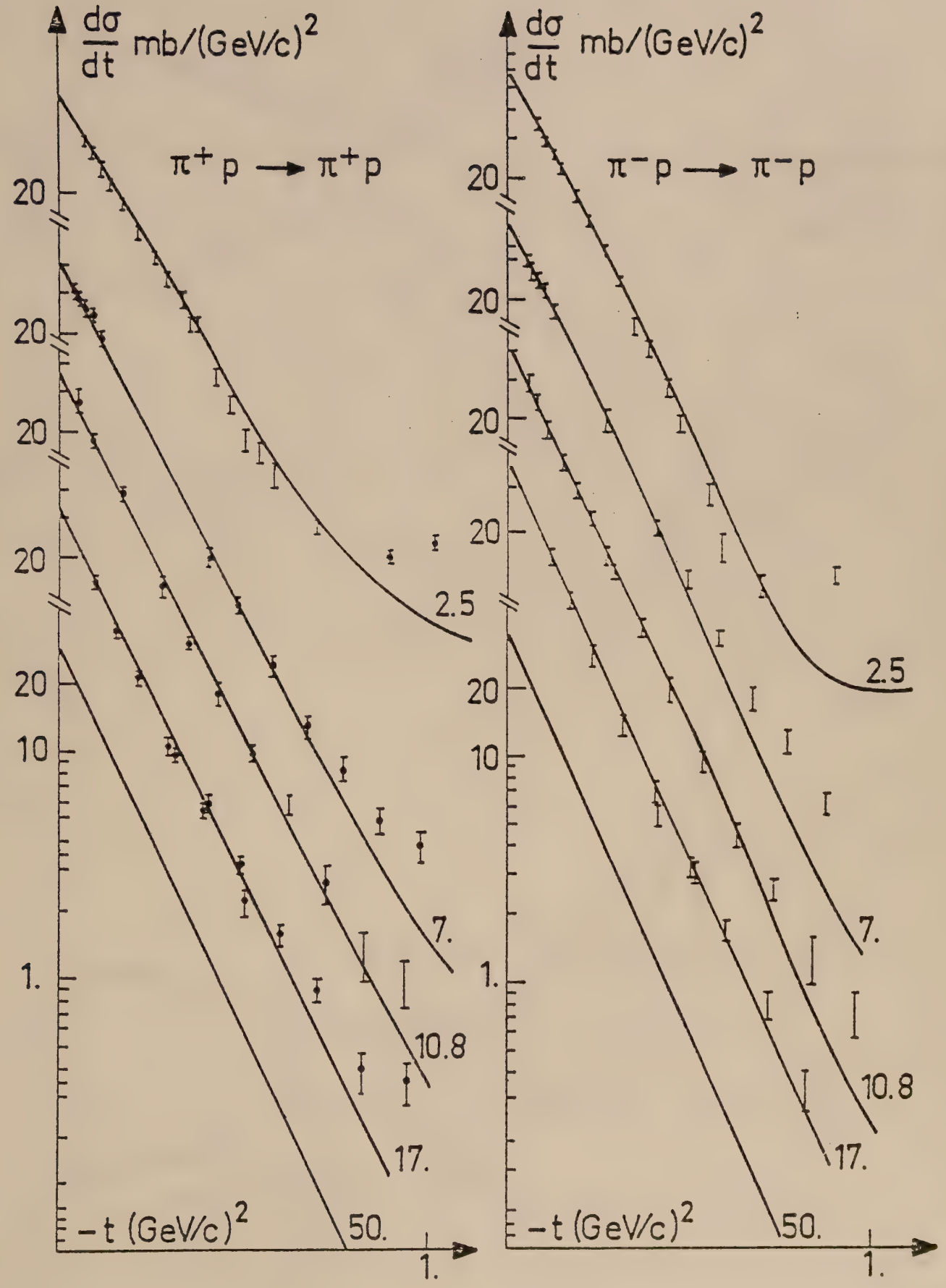
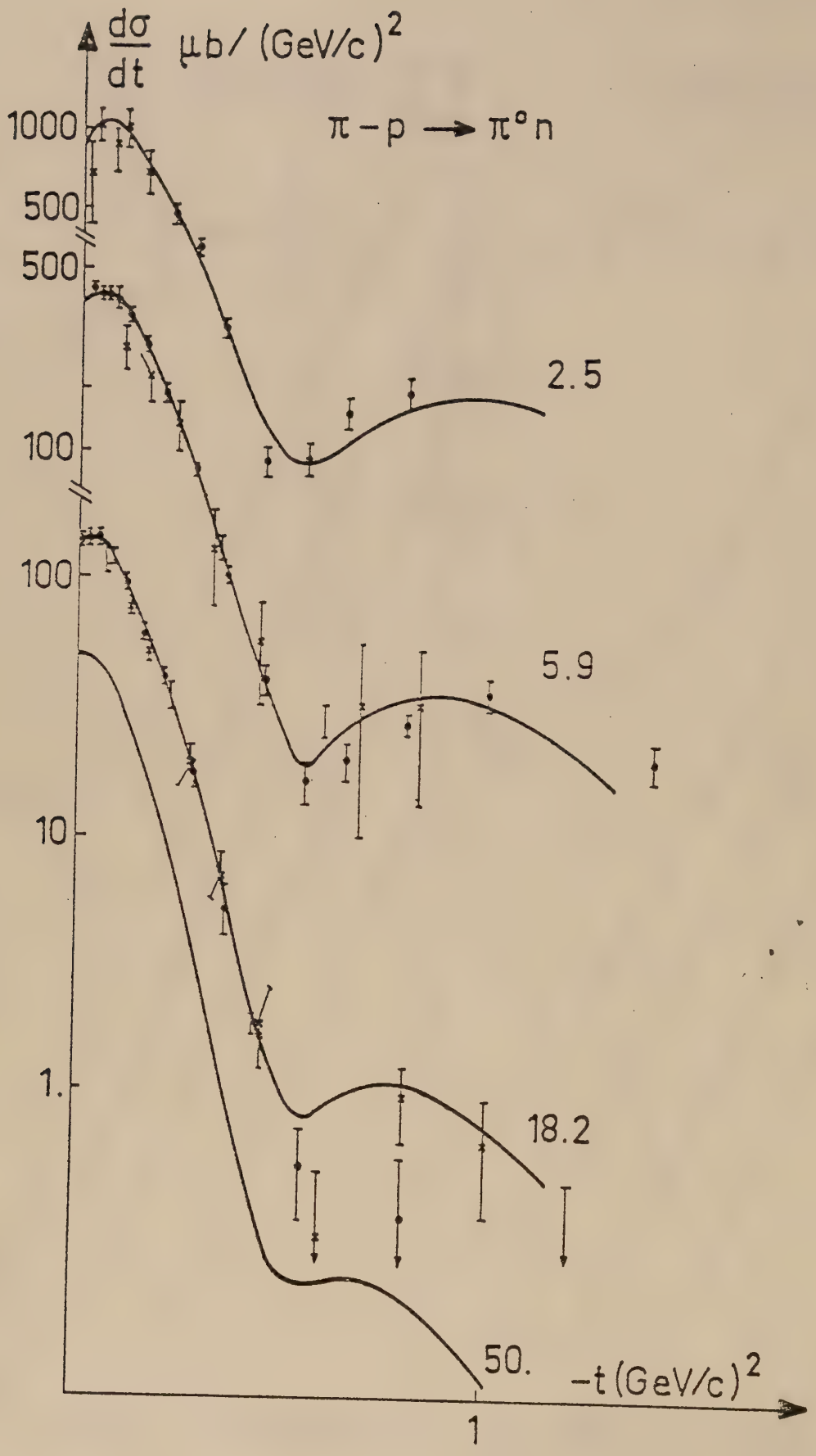

dσ/dt mb/(GeV/c)²
π⁺p → π⁺p
20
20
20
20
10
1.
2.5
7.
10.8
17.
50.
−t (GeV/c)²
1.
dσ/dt mb/(GeV/c)²
π⁻p → π⁻p
20
20
20
20
10
1.
2.5
7.
10.8
17.
50.
−t (GeV/c)²
1.

Fig.3.3


dσ/dt μb/(GeV/c)²
π⁻p → π°n
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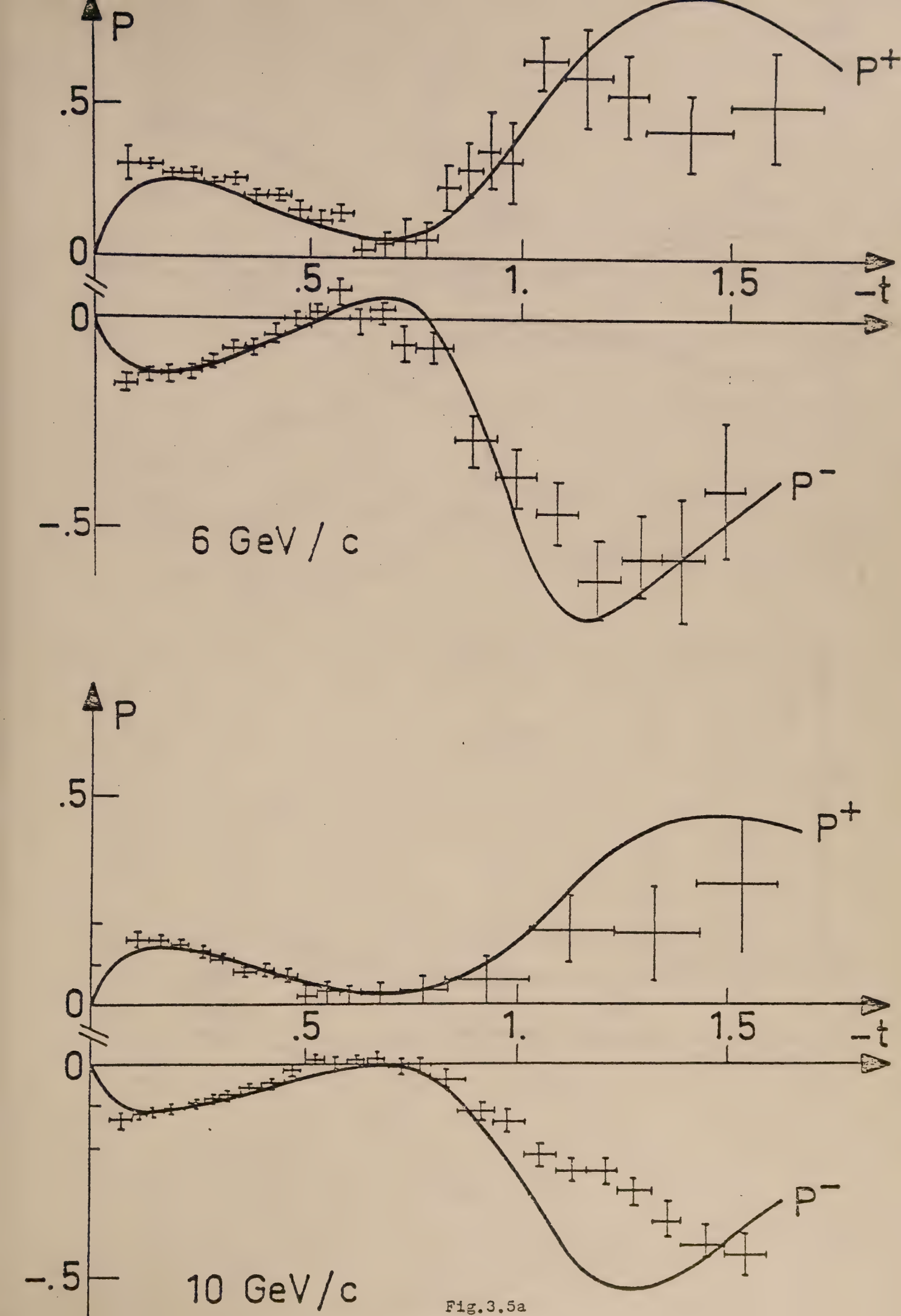
500
500
100
100
10
1.
2.5
5.9
18.2
50.
-t(GeV/c)²
1

Fig.3.4



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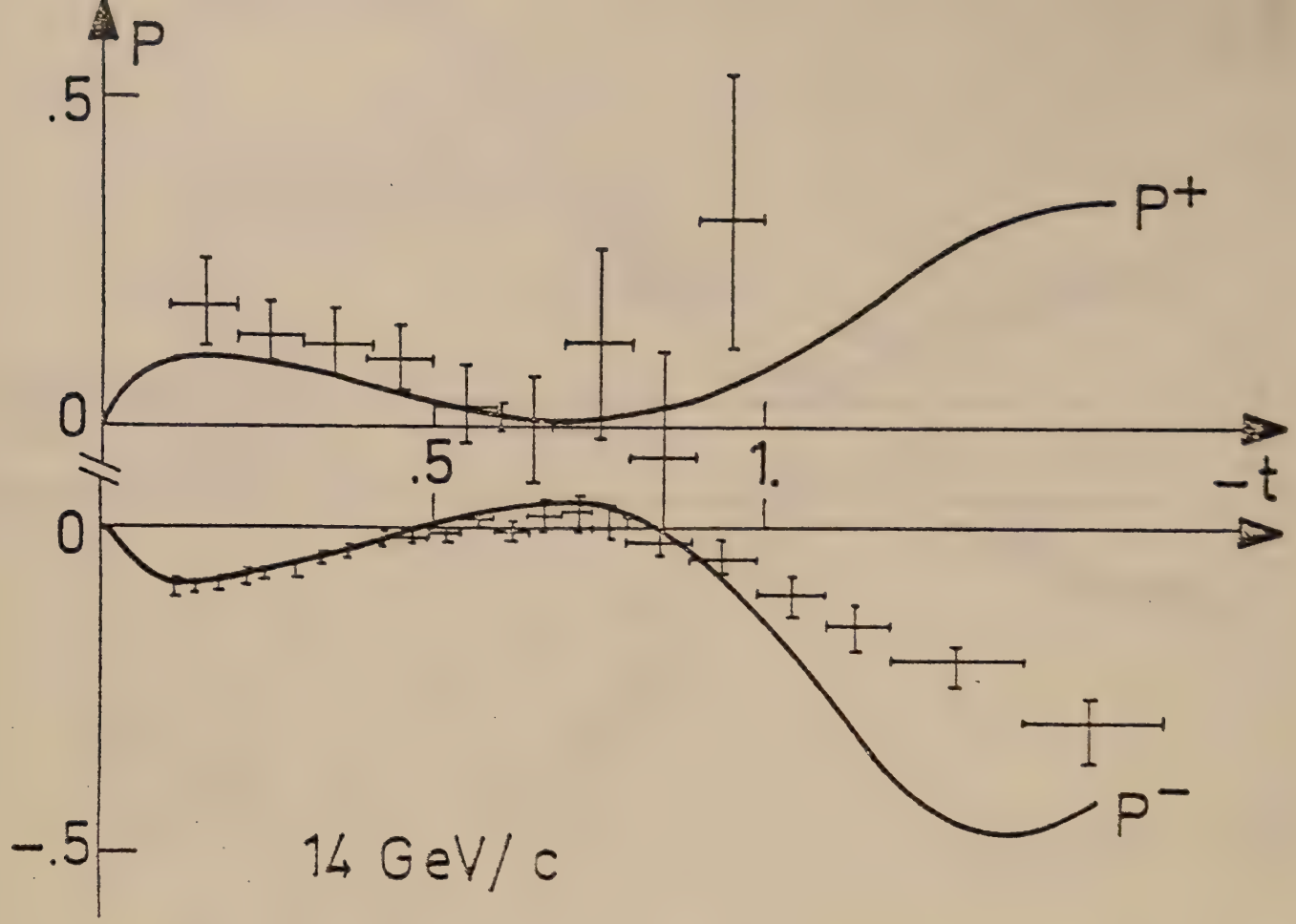

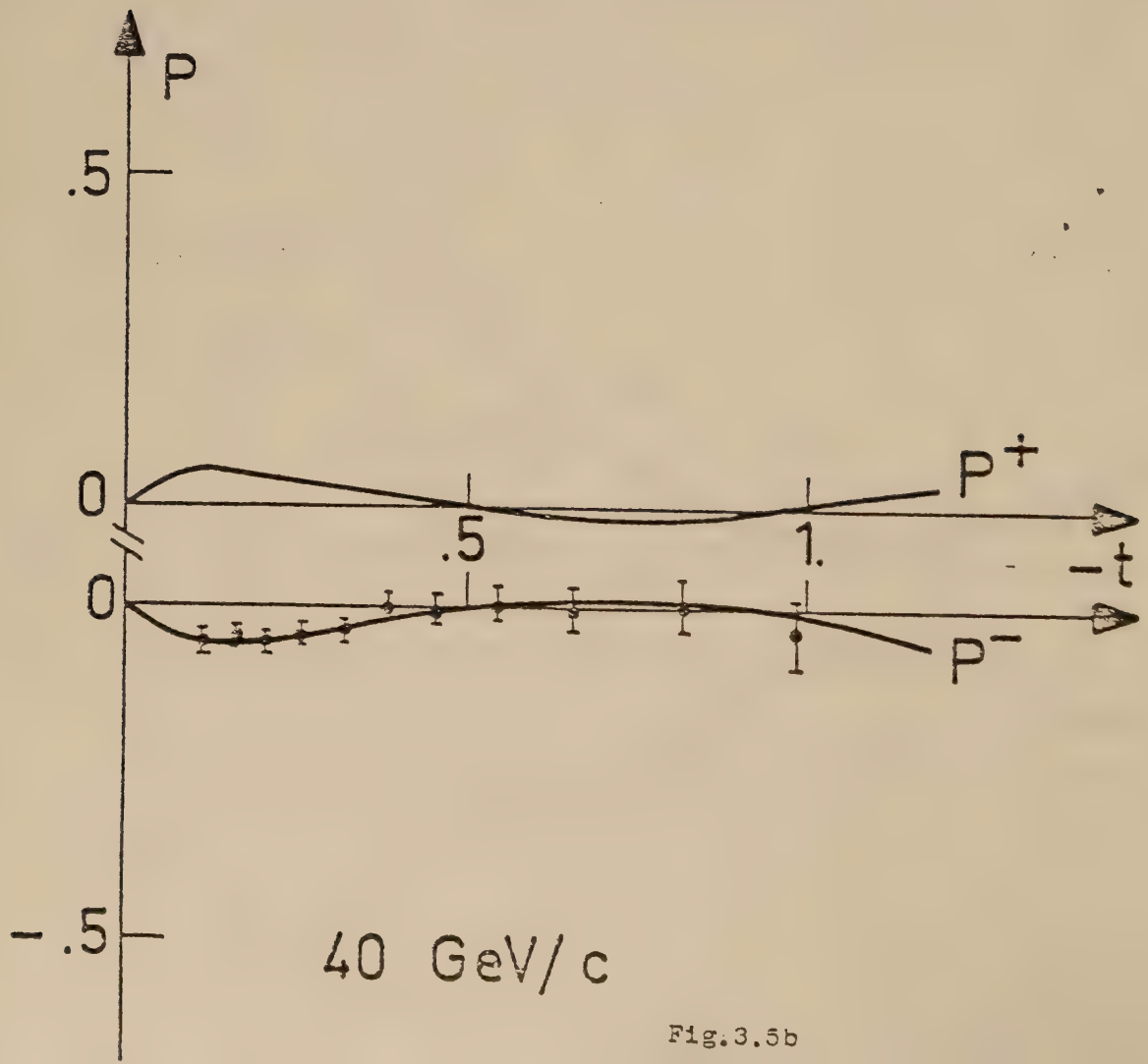

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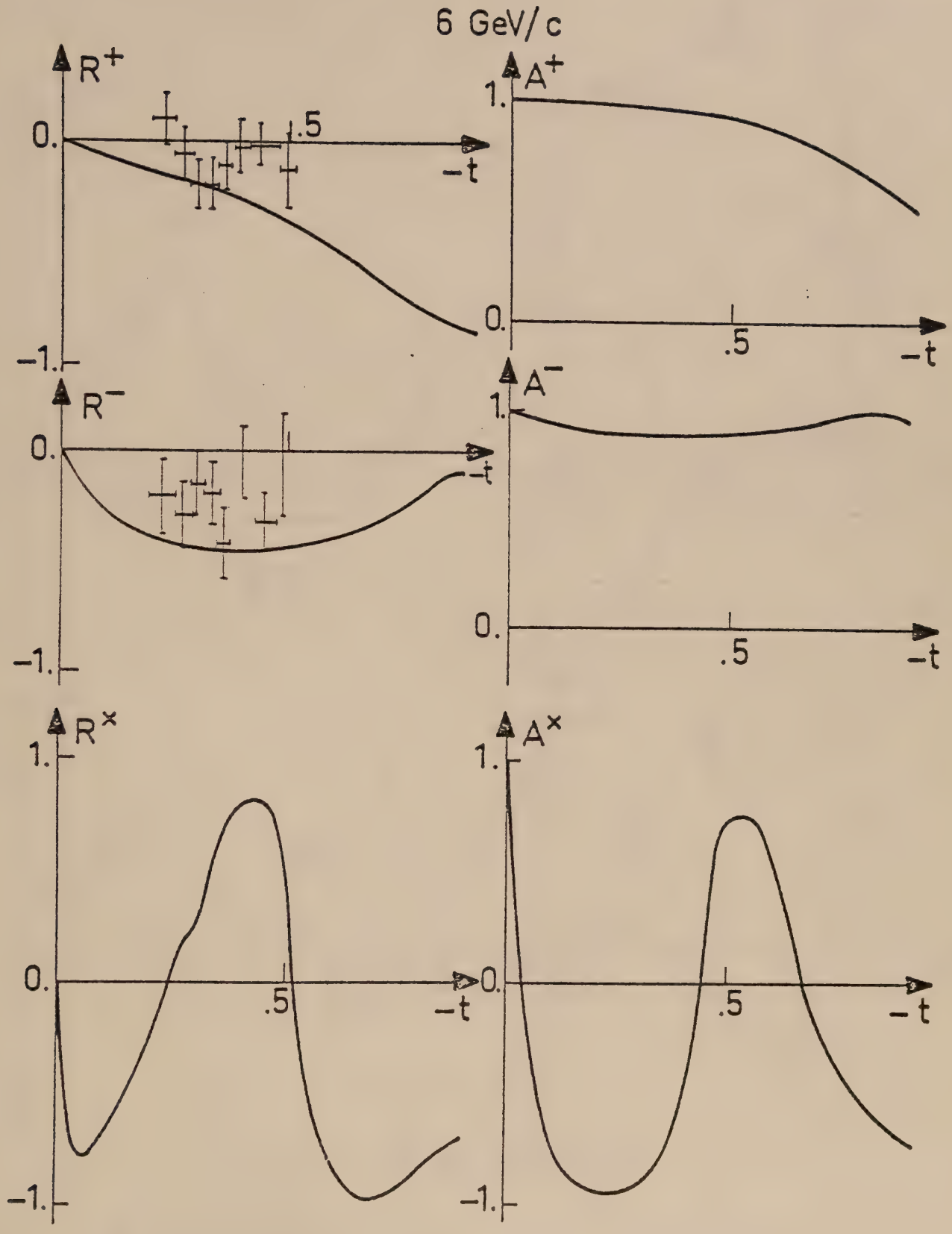

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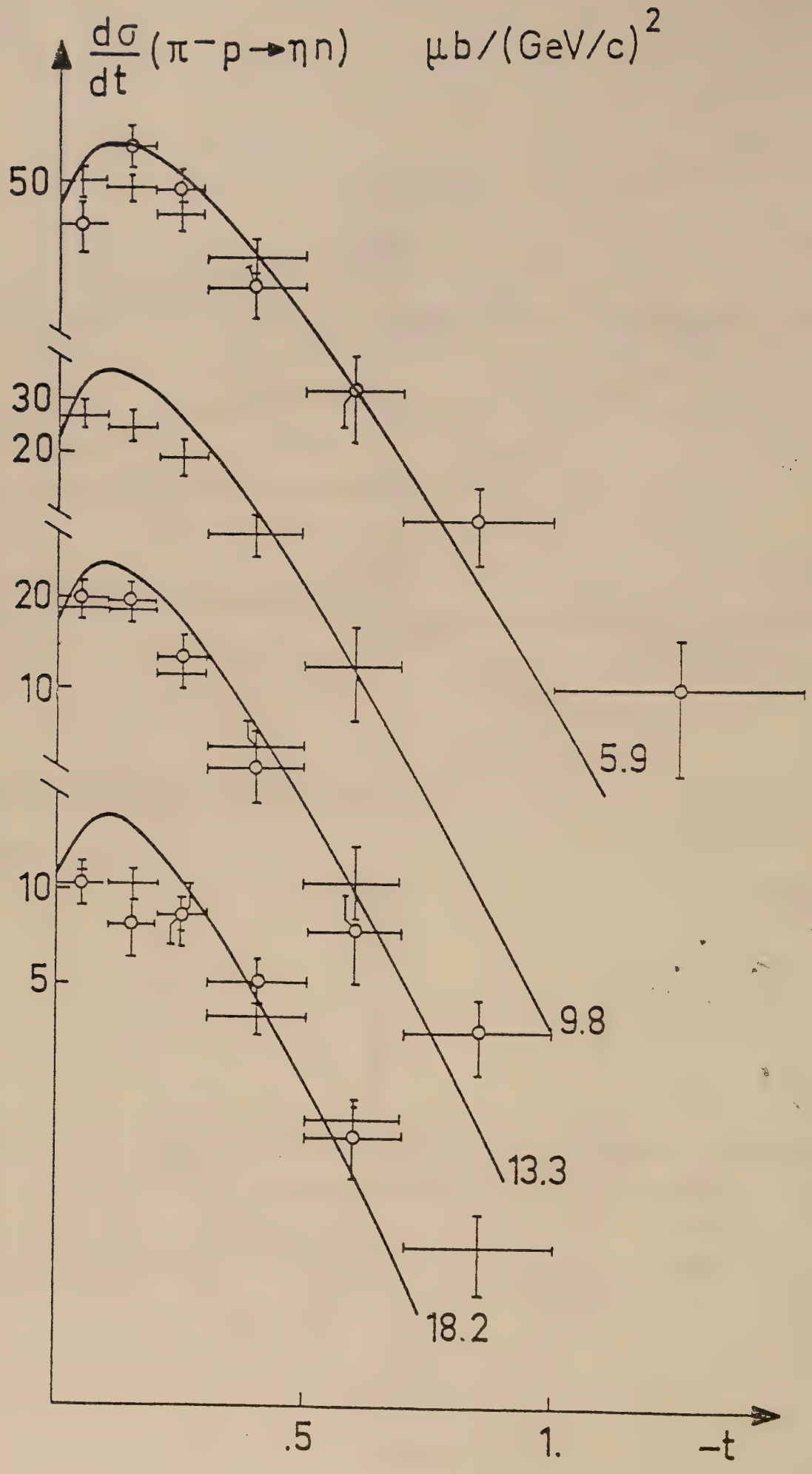

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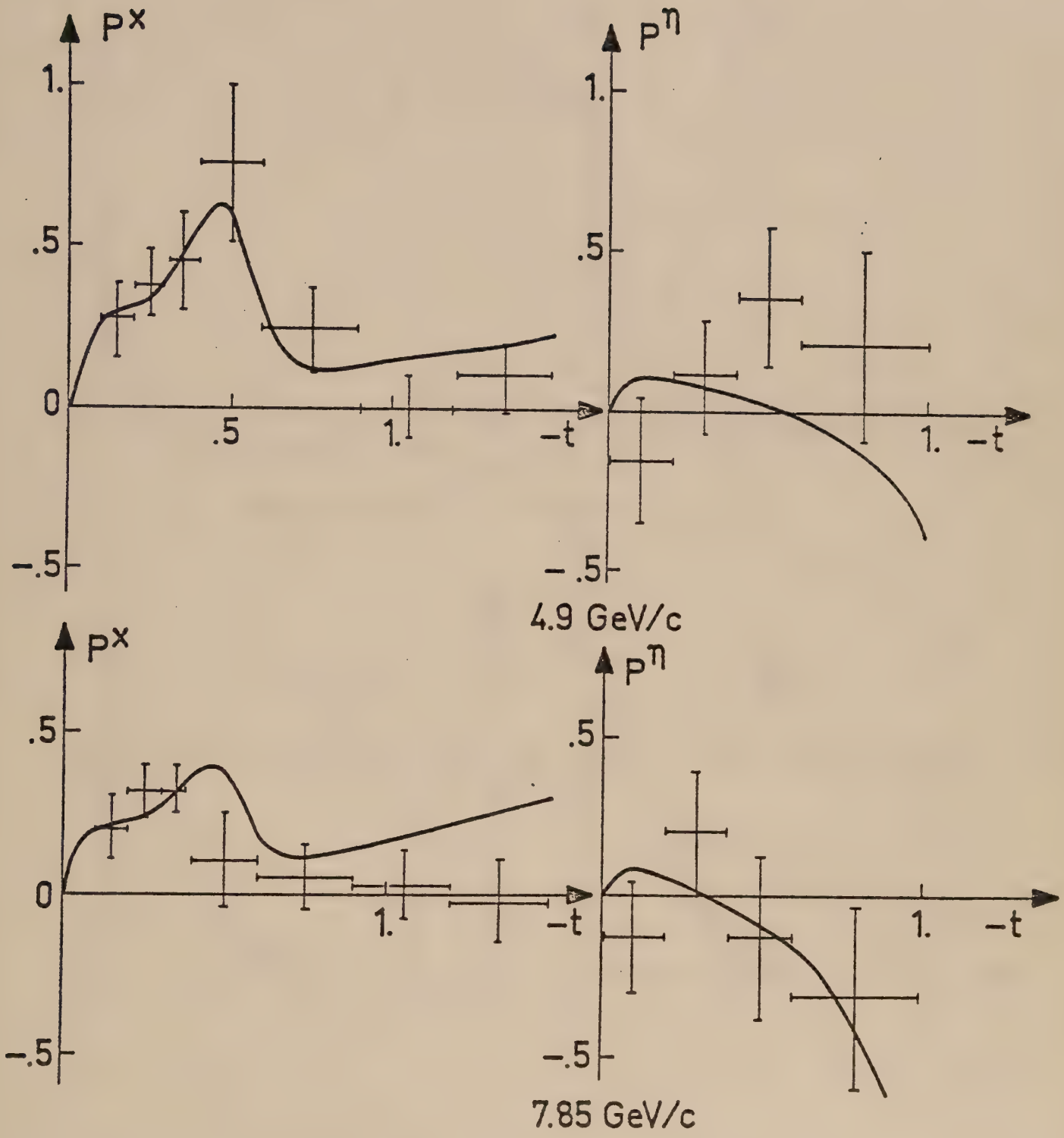

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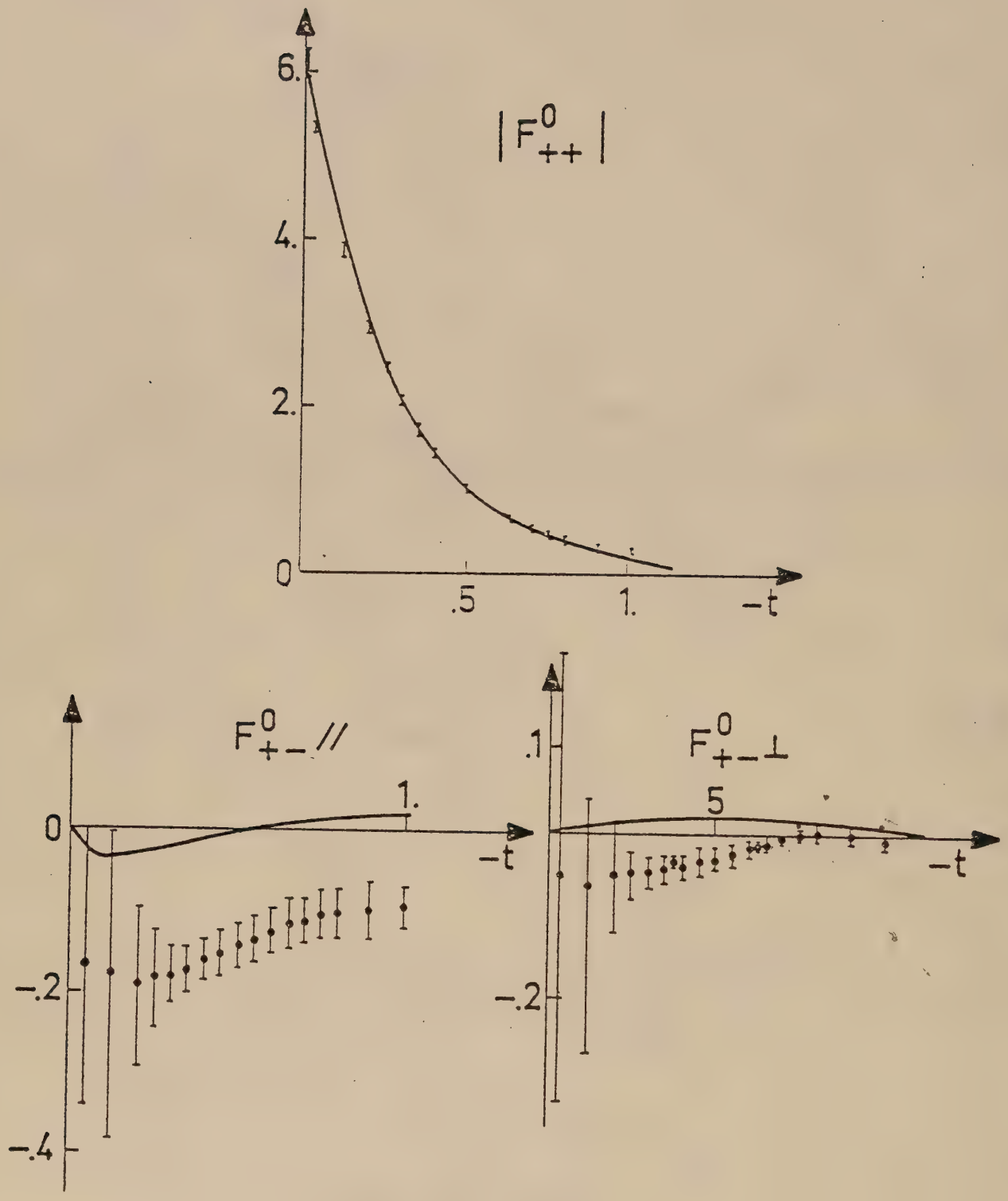

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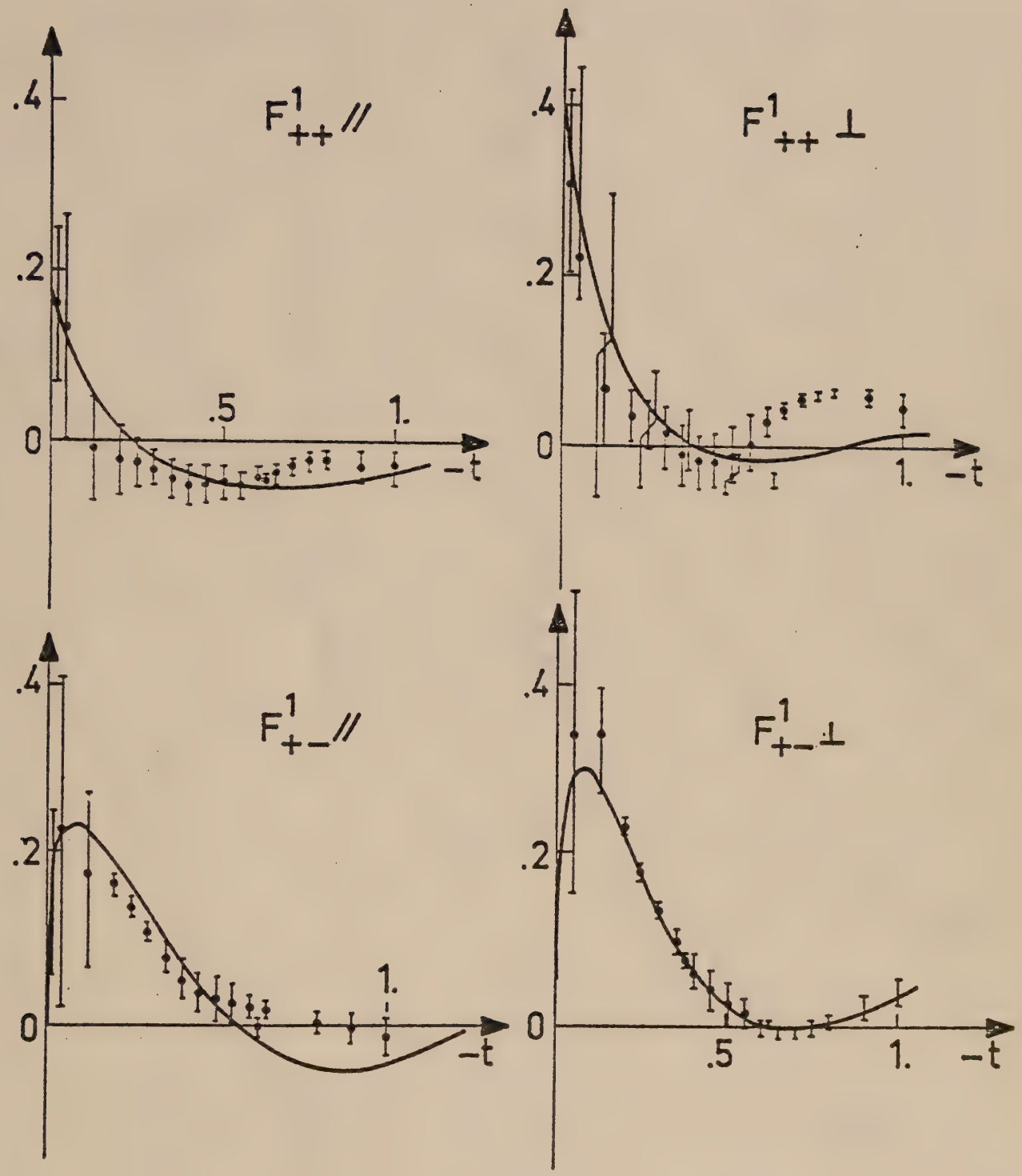

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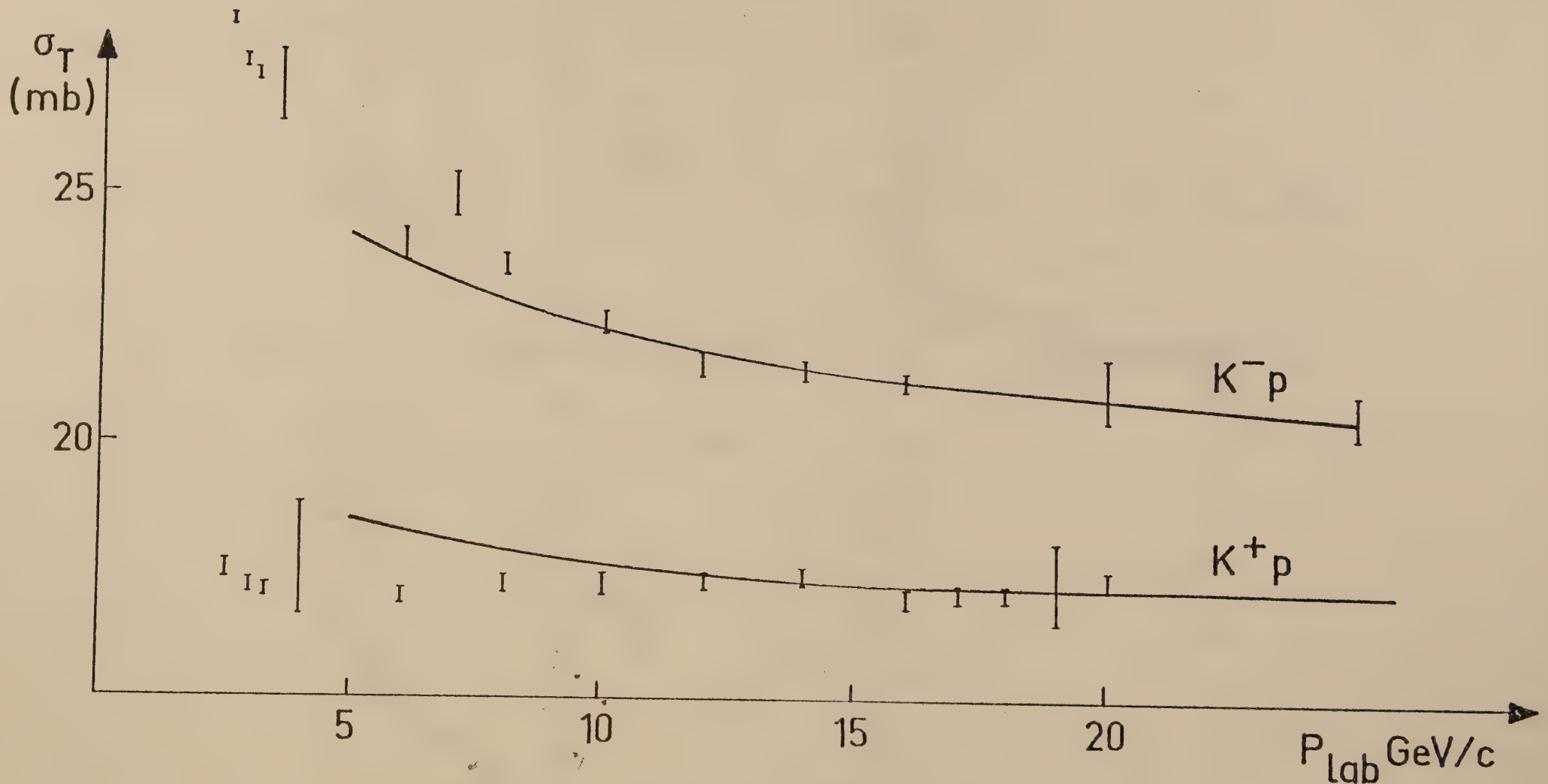

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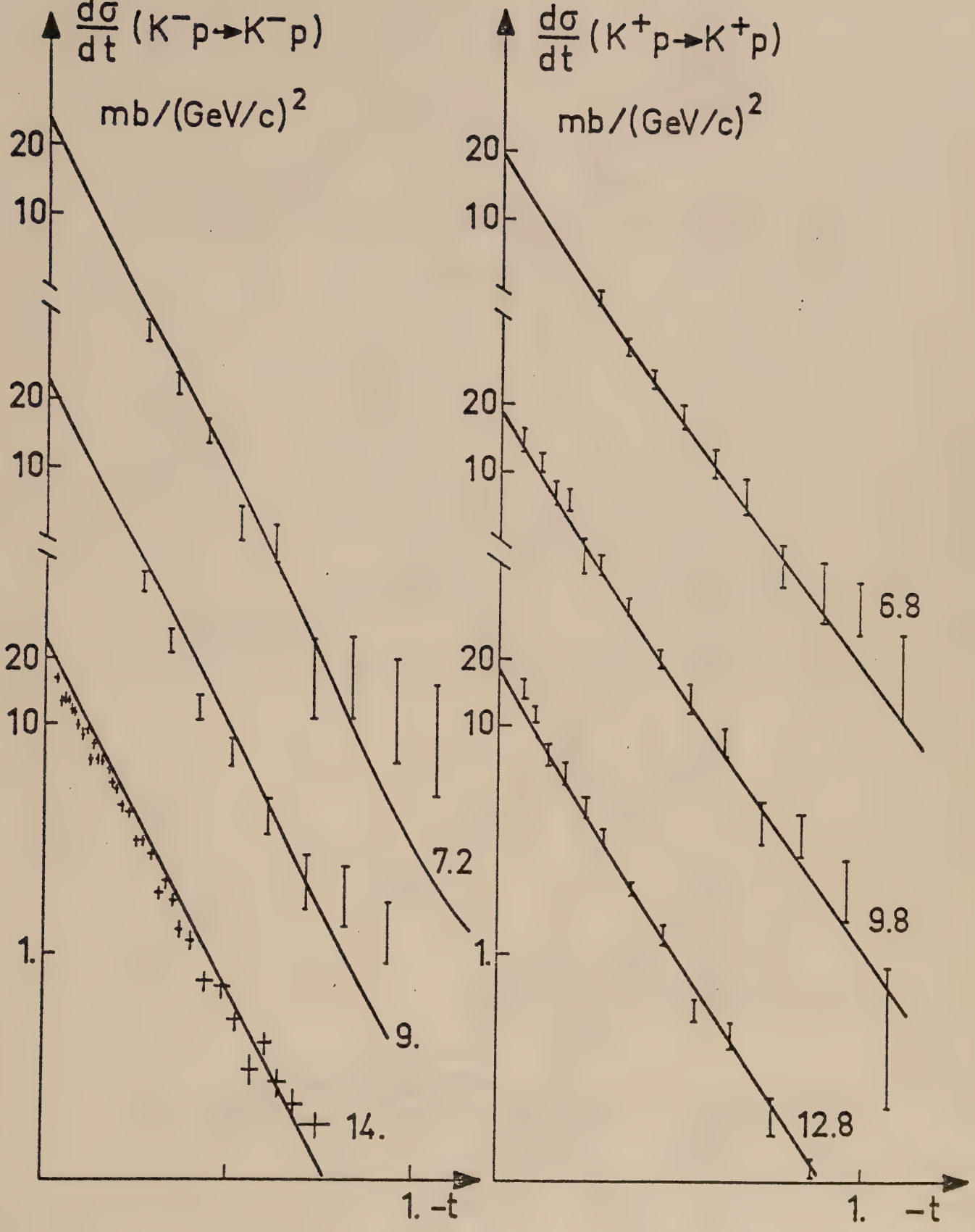
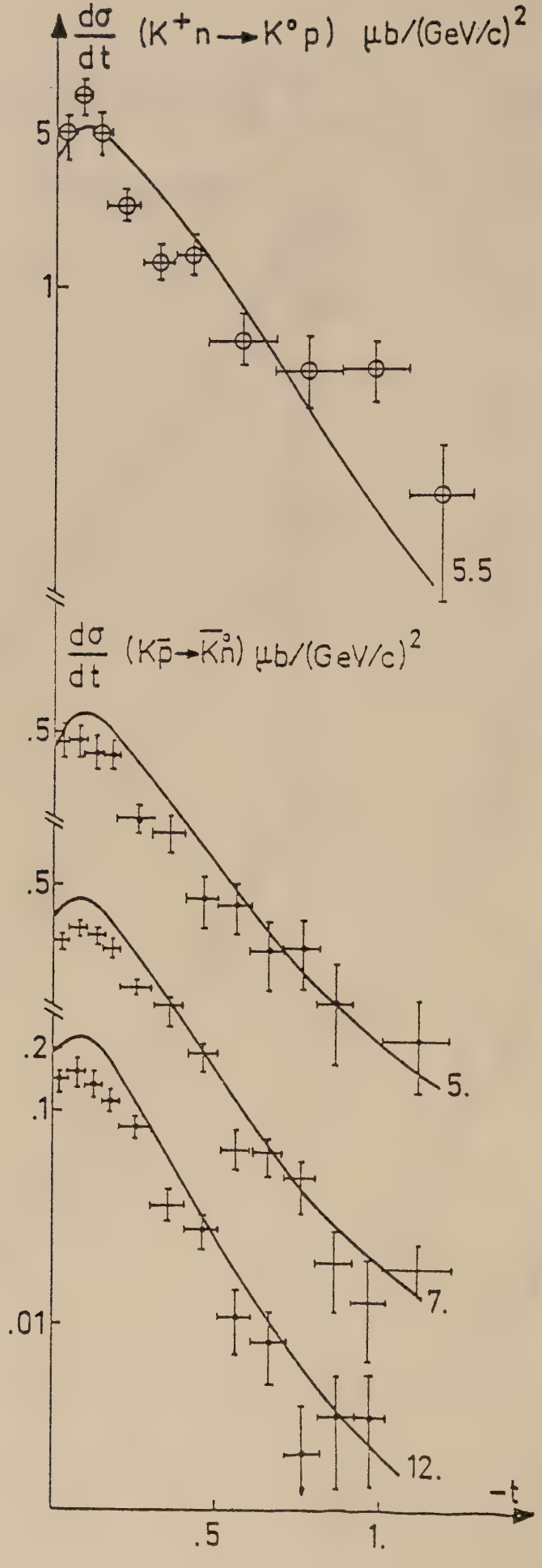

dσ/dt (K⁻p→K⁻p)
mb/(GeV/c)²
20
10
20
10
20
10
1.
7.2
9.
14.
1. -t
dσ/dt (K⁺p→K⁺p)
mb/(GeV/c)²
20
10
20
10
20
10
1.
6.8
9.8
12.8
1. -t

Fig.3.12



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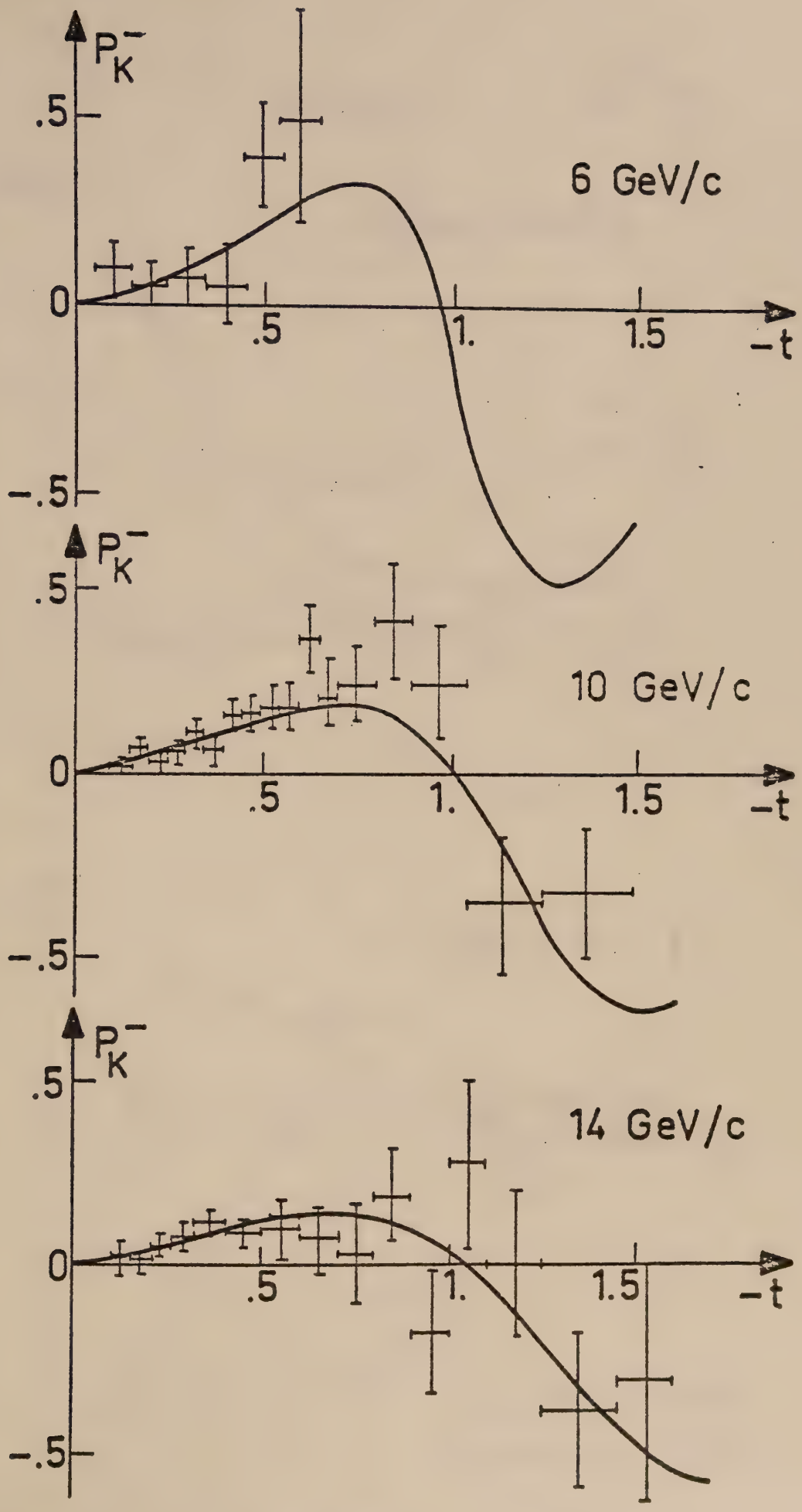


Fig.3.14a

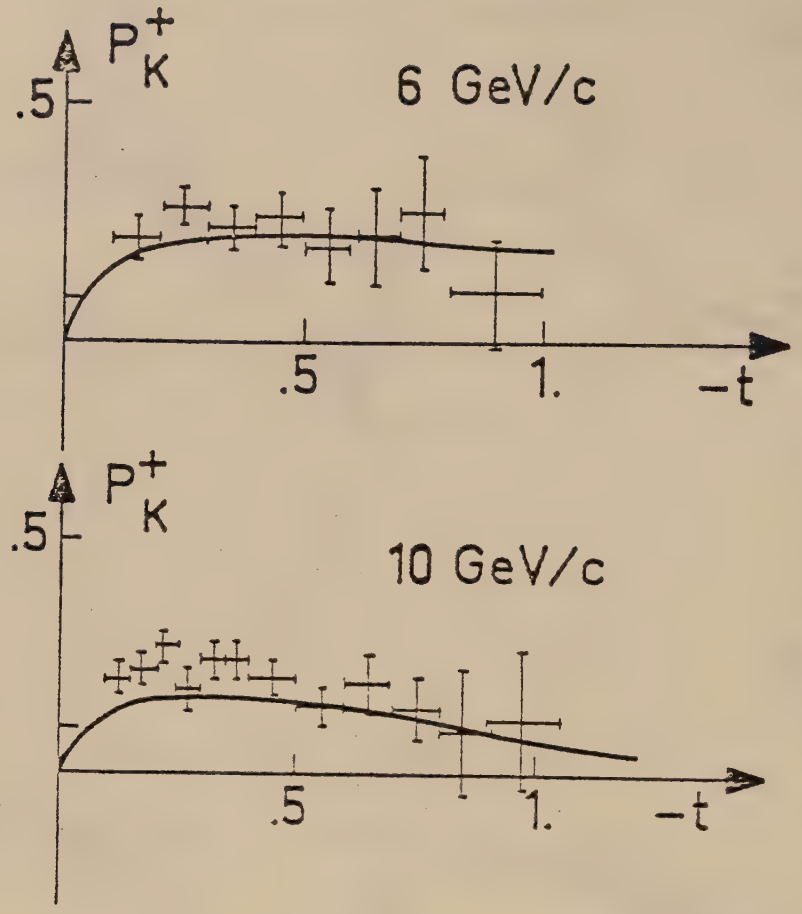
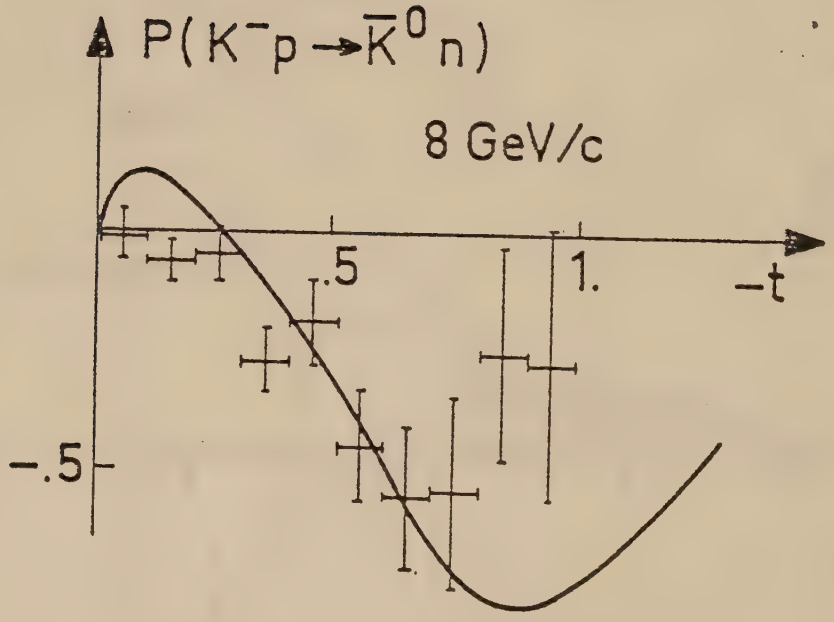


Fig.3.14b



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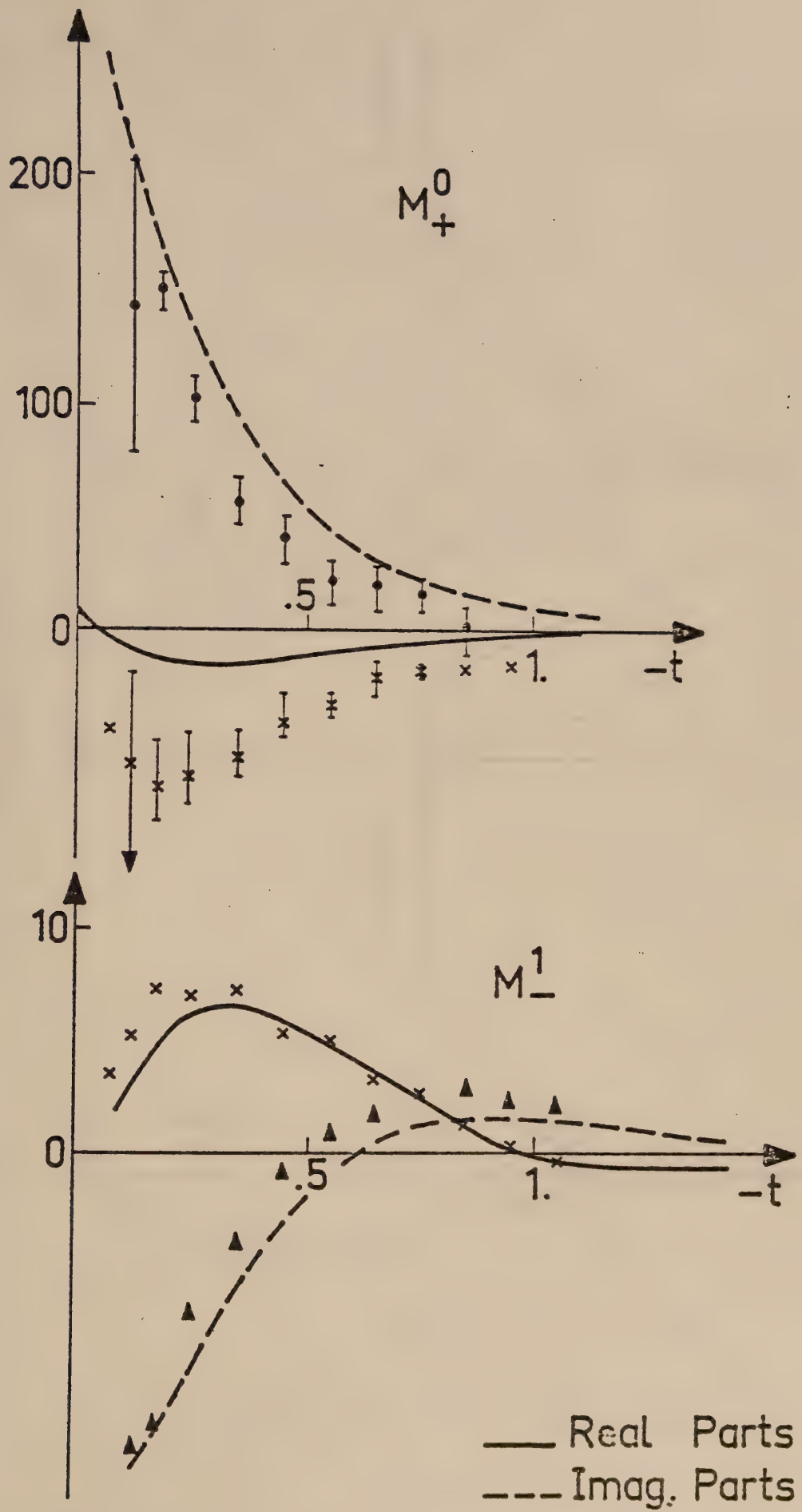

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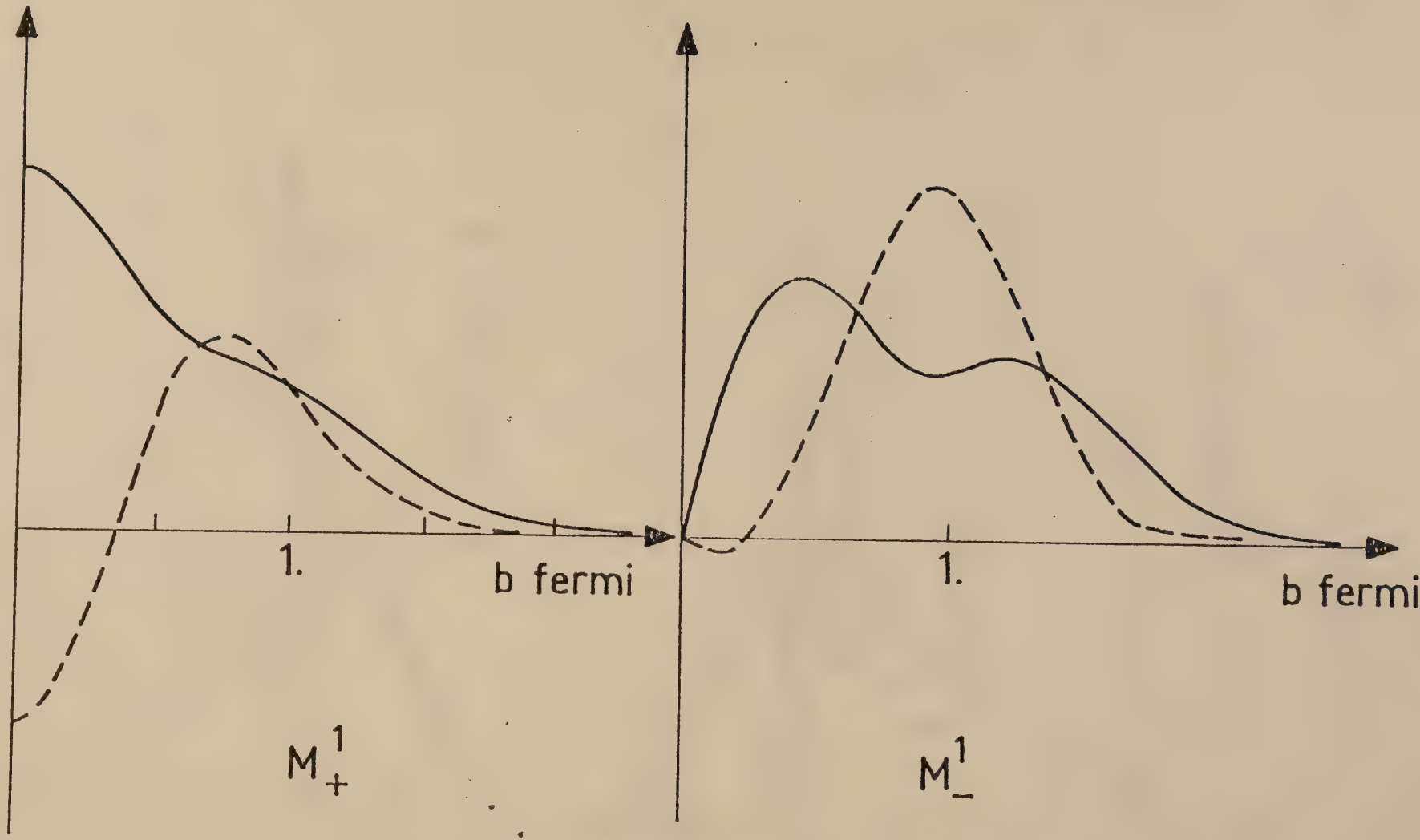
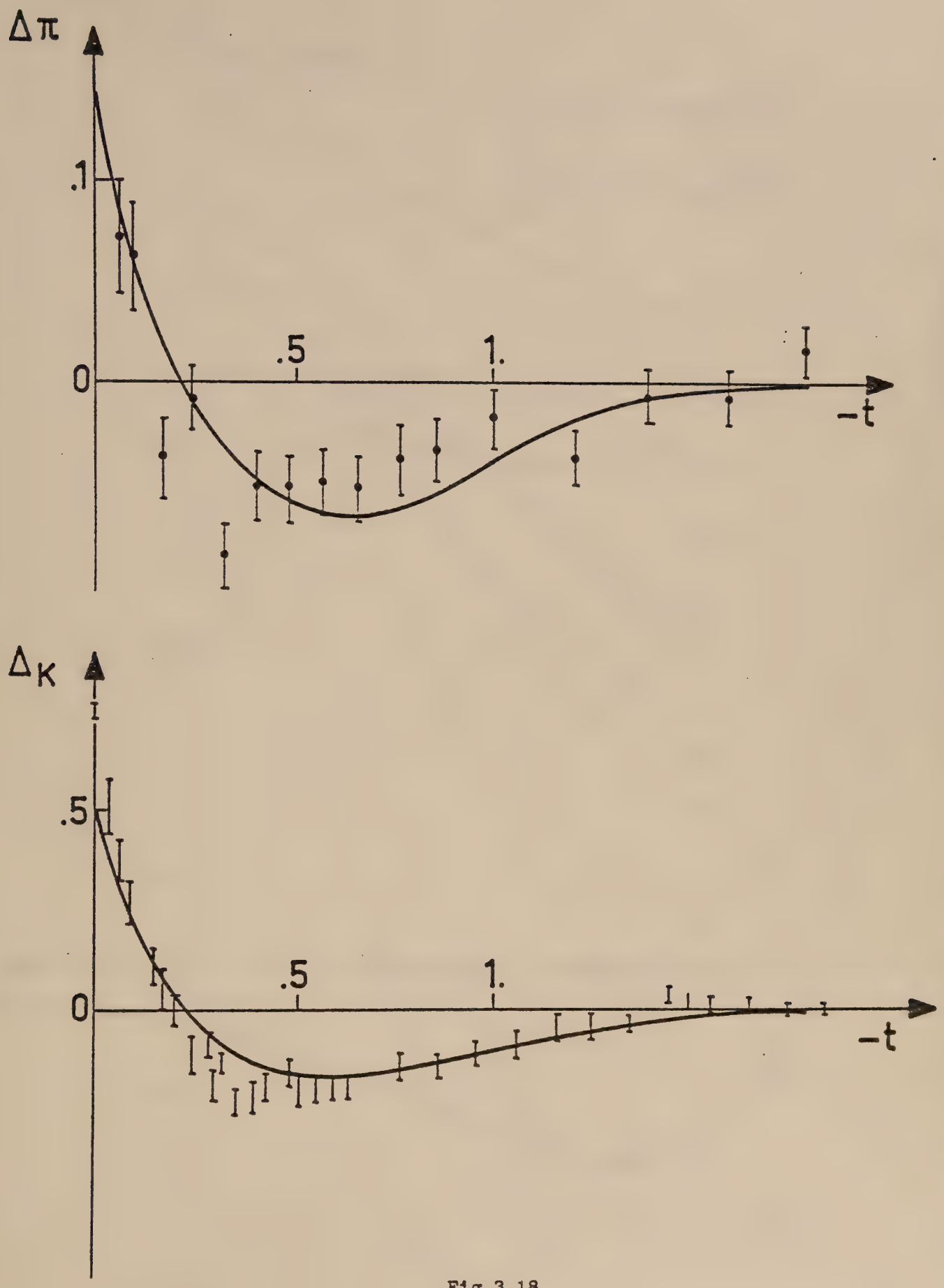


Fig.3.17



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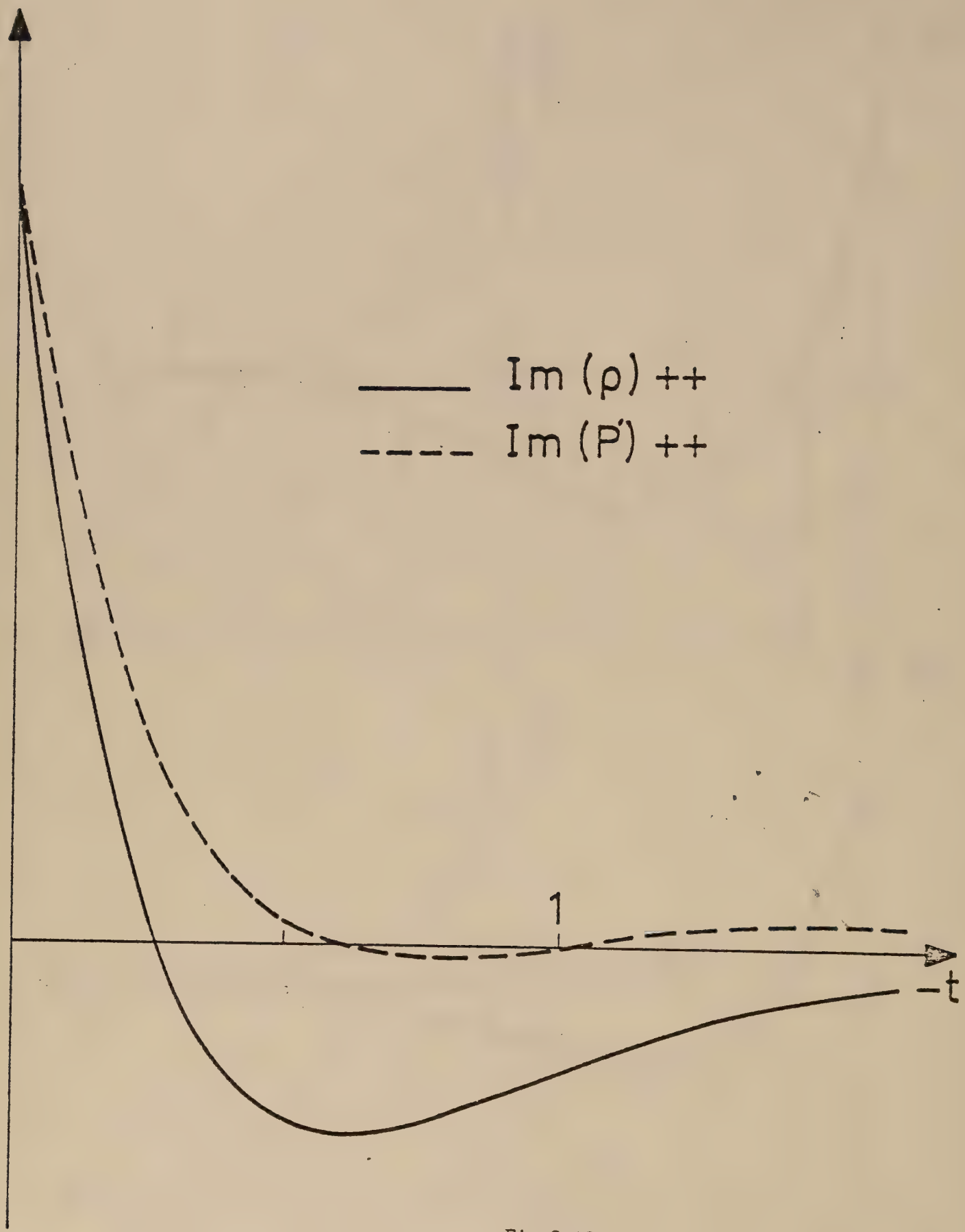


Fig.3.19

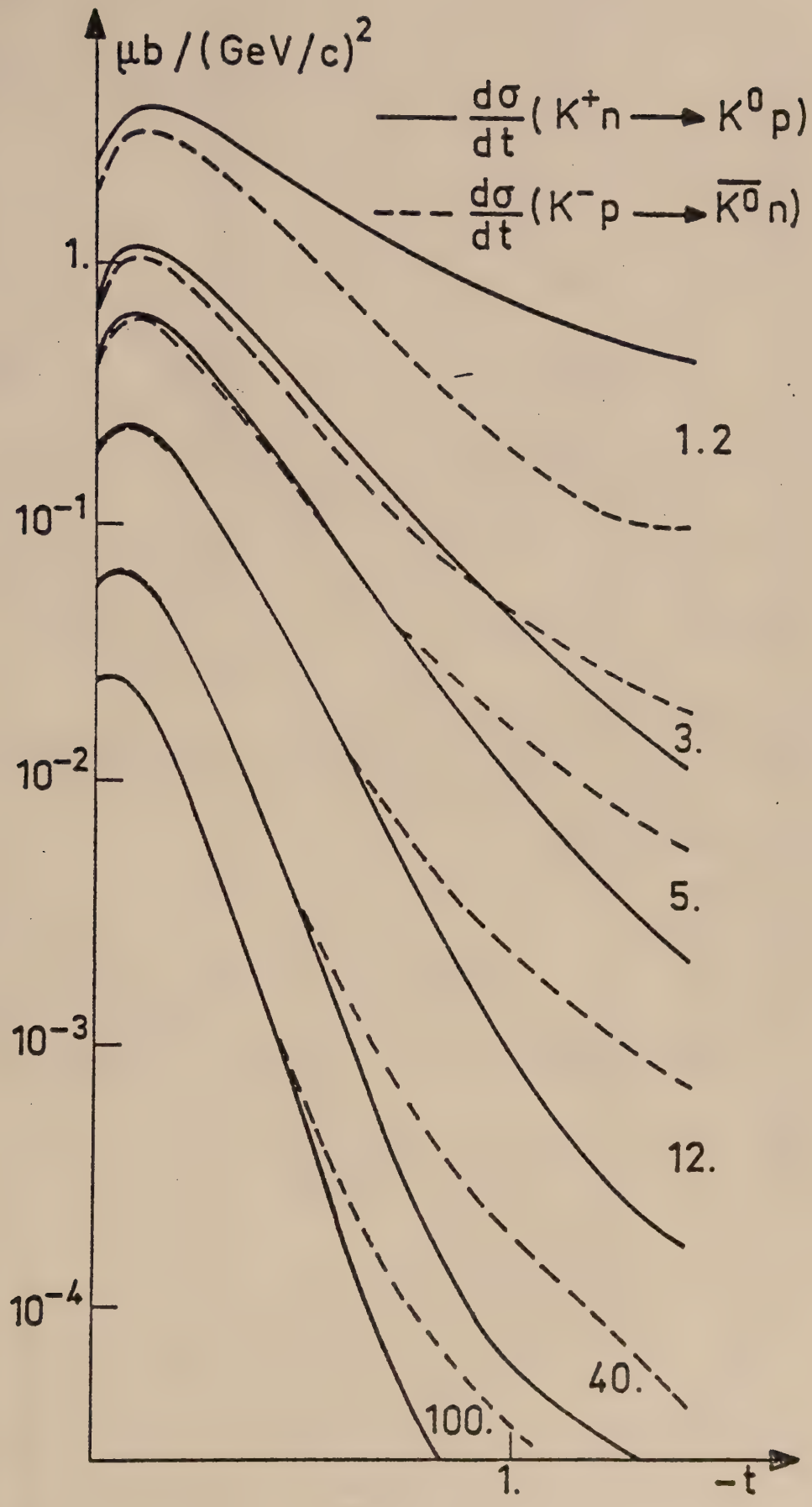
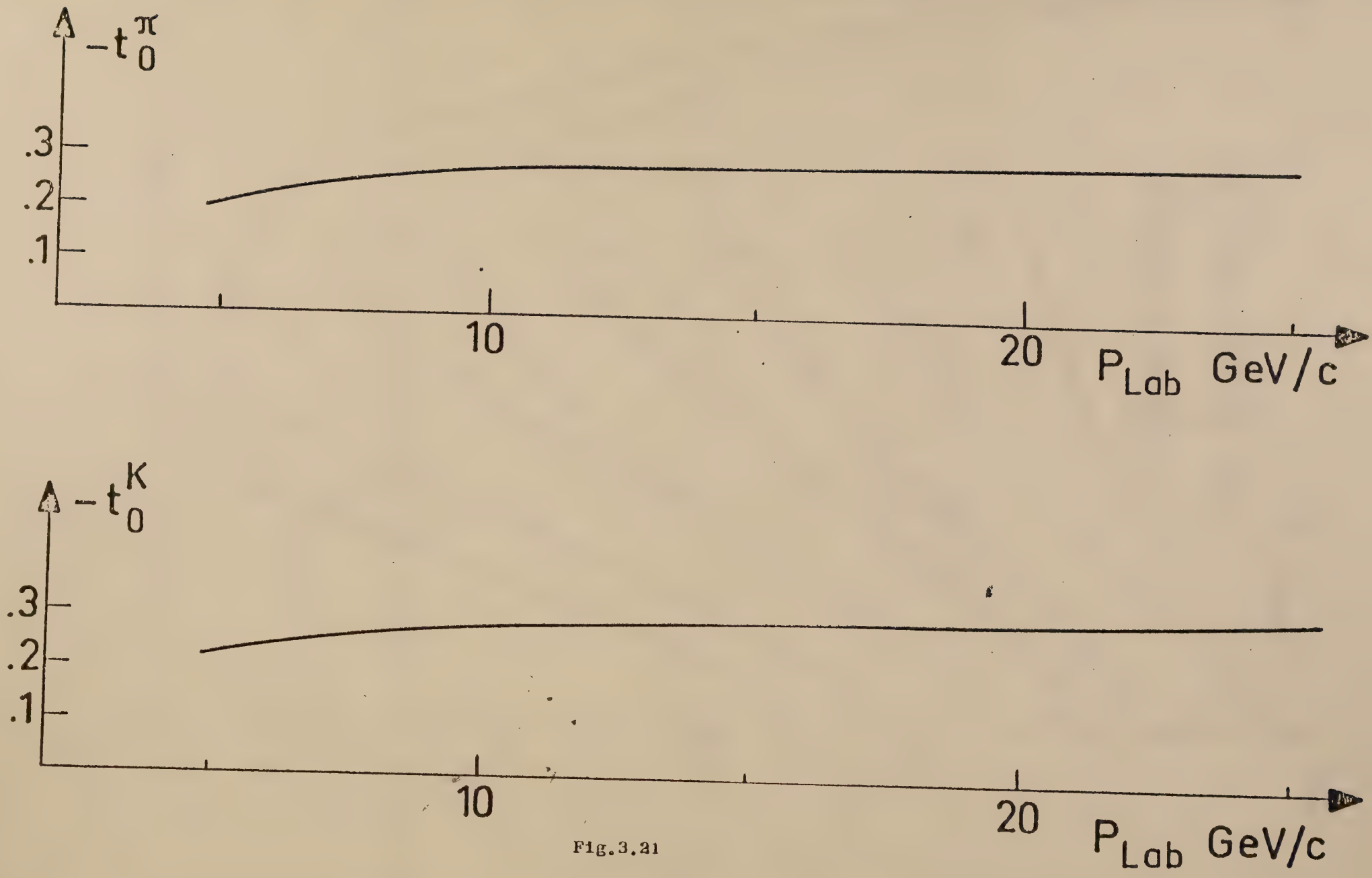

μb /(GeV/c)²
dσ/dt (K⁺n → K⁰p)
dσ/dt (K⁻p → K̄⁰n)
1.
10⁻¹
10⁻²
10⁻³
10⁻⁴
1.2
3.
5.
12.
40.
100.
1.
-t

Fig.3.20



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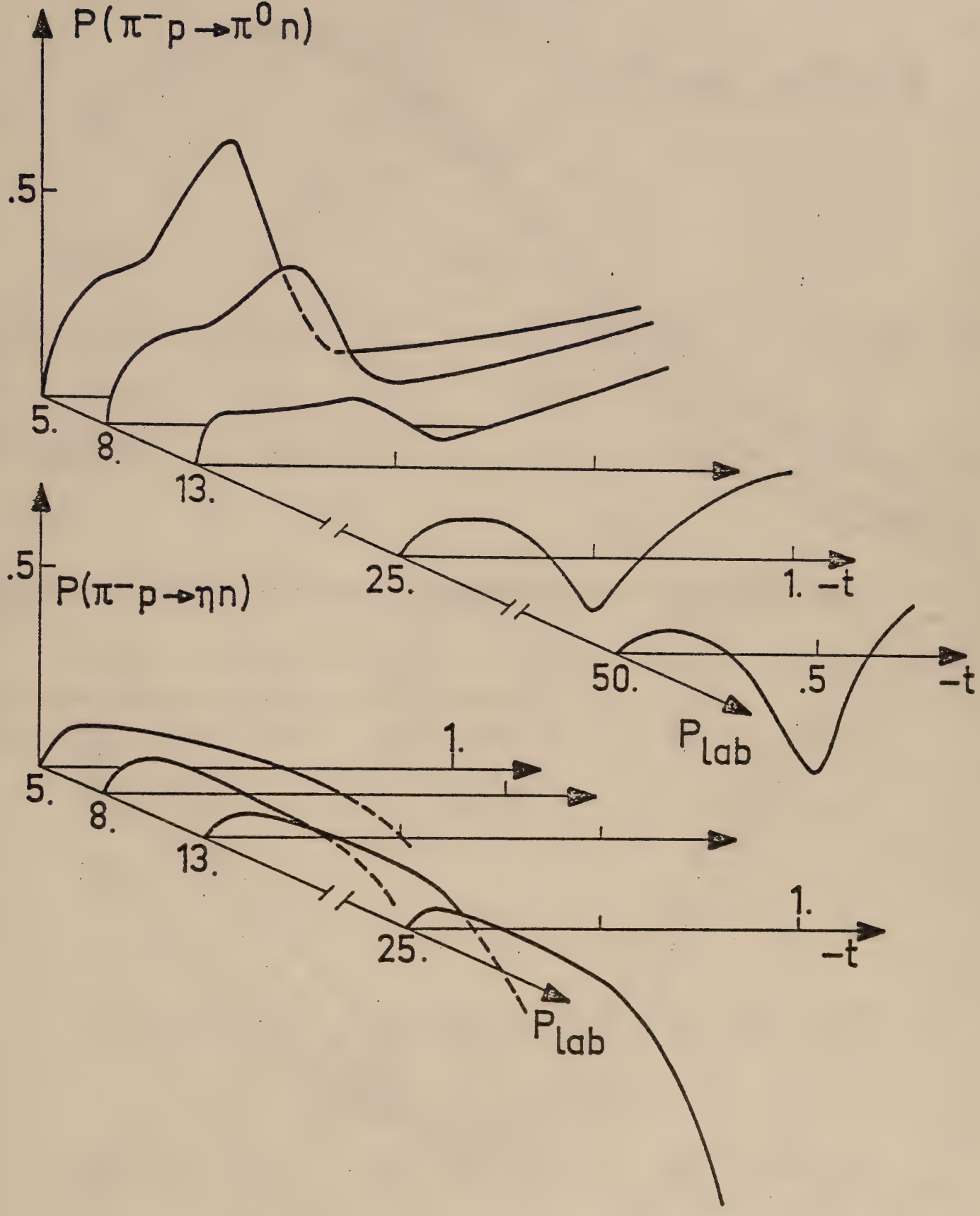


Fig.3.22

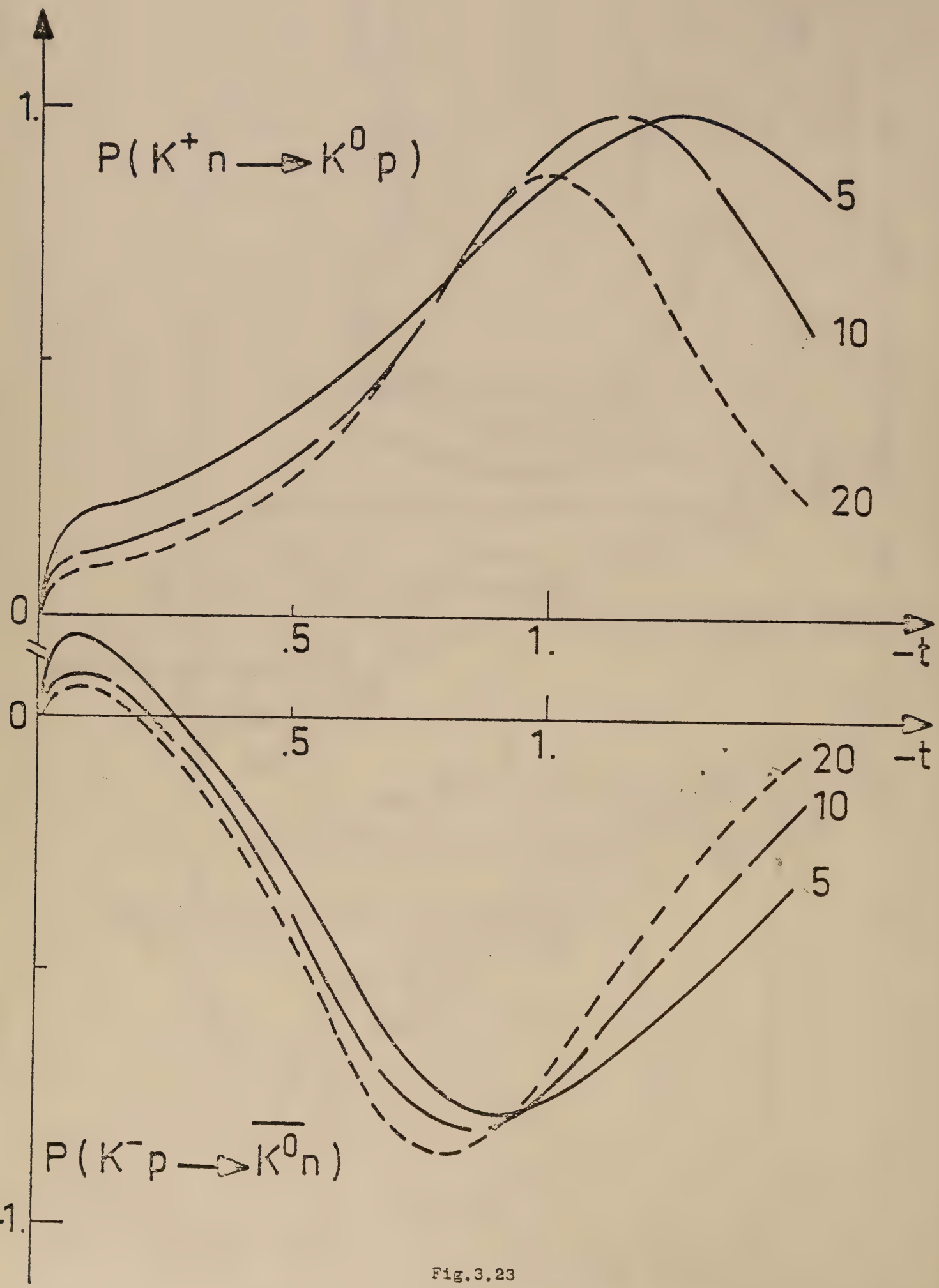


Fig.3.23

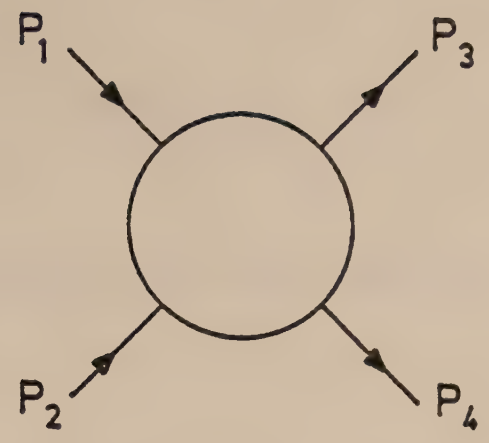


Fig.A.1

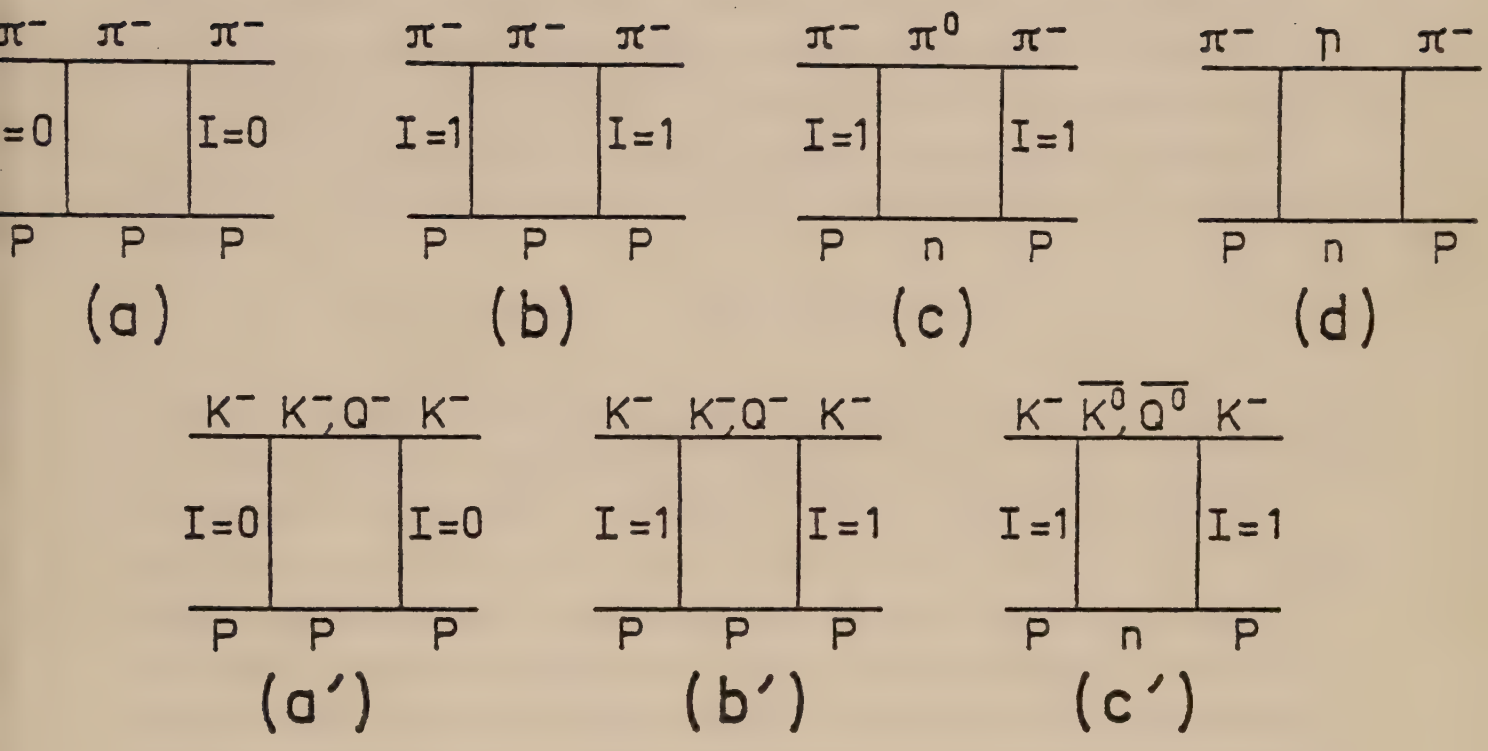


Fig.A.2

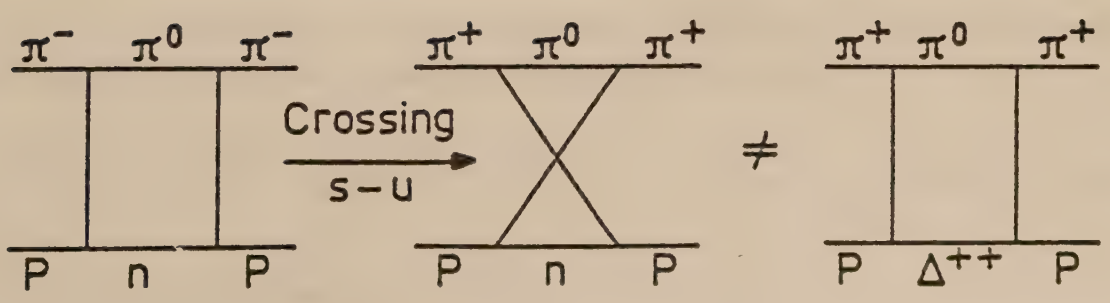


Fig.A.3

CHAPITRE III

ETUDE DES REACTIONS D'ECHANGE D'HYPERCHARGE

Nous avons vu dans le chapitre précédent comment et pourquoi nous avons construit le modèle "RPR". Avec un nombre réduit de paramètres, nous obtenons une très bonne description des réactions $\pi N \to \pi N$ et l'utilisation de SU(3) permet de prédire, sans autre paramètre, de façon satisfaisante les réactions $KN \to KN$. En particulier pour les réactions d'échange de charge $K^- p \to \bar{K}^0 n$ et $K^+ n \to K^0 p$, le modèle prédit que

$$\frac{d\sigma}{dt}(K^- p \to \bar{K}^0 n) \quad < \quad \frac{d\sigma}{dt}(K^+ n \to K^0 p) \quad , \qquad (1)$$

ce qui est observé expérimentalement[1]. Nous avons là une manifestation de la brisure de la dégénérescence d'échange. Le modèle ne rend cependant pas compte de la grandeur de l'effet, sa valeur paraissant trop faible. Ceci doit être imputé au fait qu'il n'y a pas de paramètre libre, en effet il n'y a aucune raison a priori pour que les forces des coupures dans ces réactions soient comparables à celles qui sont présentes dans πN. On peut penser qu'un calcul plus réaliste, tenant compte des différents états intermédiaires donnerait une meilleure estimation de la brisure de la dégénérescence d'échange. Dans le présent chapitre, nous allons nous attacher à la description des réactions d'échange d'hypercharge : $\pi N \to K\Sigma/\Lambda$ et $\bar{K}N \to \pi\Sigma/\Lambda$ qui sont intéressantes à plus d'un titre.

(a) La brisure de la dégénérescence d'échange est plus importante que dans les réactions d'échange de charge KN, ainsi

$$\frac{d\sigma}{dt}(K^- p \to \pi^- \Sigma^+) \,/\, \frac{d\sigma}{dt}(\pi^+ p \to K^+ \Sigma^+) \approx 1.5 \quad \text{à} \quad p_{lab} = 6 \text{ GeV/c} \; . \qquad (2)$$

(b) A la différence des réactions d'échange de charge où les amplitudes d'hélicité flip (n=1) dominent, ce sont les composantes (n=0)

qui sont importantes. Ces amplitudes sont intéressantes car elles n'ont pas un comportement de pôle de Regge simple. On peut espérer avoir ici des informations claires sur leur structure en transfert.

Des analyses en amplitudes[2-7] des réactions méson-baryon suggèrent notamment que les amplitudes d'échange de vecteurs (ρ,ω,K_V) ou de tenseurs ($A_2 f_o, K_T$) ont des comportements universels c'est-à-dire ne dépendant de la réaction que par leurs couplages. Par exemple dans les échanges d'hypercharge, on peut remplacer respectivement le K_V et le K_T par le ρ et le A_2, en tenant compte des différences de couplages et de trajectoires entre les K^* et (ρ,A_2)[7] et obtenir une bonne description des résultats expérimentaux. Il s'agit ici non du terme de pôle seul mais de toute l'amplitude d'échange portant dans la voie t les nombres quantiques du vecteur et du tenseur échangés comme pôles de Regge.

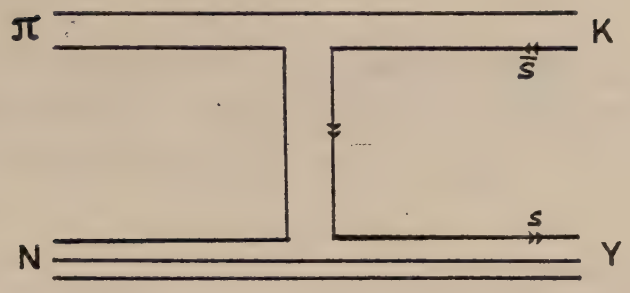
(c) Il existe également des mesures de $d\sigma/dt(t=0)$ pour les réactions $\pi^- p \to K^+\Sigma^-$ et $K^- p \to \pi^+\Sigma^-$ [8], qui ont une voie t exotique (échange d'hypercharge et double échange de charge). Dans ce cas, les seules contributions - dans un modèle d'échange - viennent de coupures à au moins deux Reggeons. Il est intéressant de voir ce que donne le modèle dans ce cas précis.

Nous avons choisi ici de ne laisser aucun paramètre libre et de tout calculer a priori afin de voir ce que donne le modèle. Dans ce qui suit nous présentons de façon très détaillée et sans doute fastidieuse le modèle et ses contributions ainsi que le calcul des paramètres.

LE MODELE RPR ET LES REACTIONS D'ECHANGE D'HYPERCHARGE

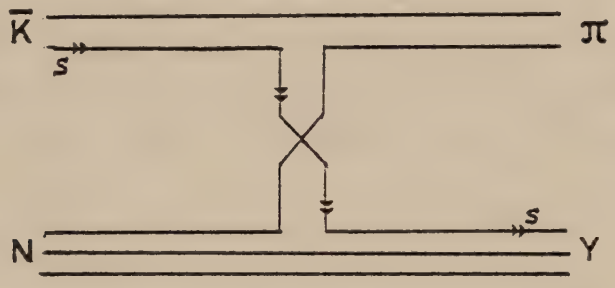
A - Les termes de pôles

Dans les réactions d'échange d'hypercharge $\pi N \to KY$ et $\bar{K}N \to \pi Y$ (Y est une notation pour Σ ou Λ) nous avons 2 types d'échanges possibles : un échange de type vecteur $K^*(892)\ J^P = 1^-$ (K_V) et un échange de type tenseur $K^{**}(1420)\ J^P = 2^+$ (K_T) qui sont de signature opposée. Nous représentons les réactions à l'aide de diagrammes de quarks. Ainsi $\pi N \to KY$ est représentée par un diagramme plan qui correspond à la partie $e^{-i\pi\alpha}$



$e^{-i\pi\alpha}$

du pôle de Regge. Ceci indique que la voie $\pi N \to KY$ est non exotique c'est-à-dire que $\langle \text{Im (Résonances)} \rangle \neq 0$, puisque le terme de Regge contient de la partie imaginaire. En ce qui concerne la réaction $\bar{K}N \to \pi Y$, le diagramme de quarks associé est non plan



1

Un tel diagramme correspond à la partie 1 du pôle de Regge, on a par dualité $\langle \text{Im (Résonances)} \rangle = 0$. Les diagrammes de quarks donnent une extension de la notion d'exotisme à savoir sont exotiques non seulement les voies où les nombres quantiques sont exotiques au sens de SU(3) mais également celles où le diagramme de quarks est non plan. Ainsi la voie $\bar{K}N \to \pi Y$ contient, dans la voie s , les résonances Y^* mais on a pu montrer que la valeur moyenne de leurs contributions à la partie imaginaire de l'amplitude est nulle. Afin d'être cohérents avec les notations des appendices A et B nous noterons pour les pôles de Regge vecteurs et tenseurs :

$$V = -g(r-\Phi) \qquad T = g(r+\Phi) \quad , \tag{3}$$

où r et Φ sont respectivement associés à la partie 1 et à la partie $e^{-i\pi\alpha(t)}$ c'est-à-dire :

$$r = \Gamma(1-\alpha(t))\ (\nu/\nu_o)^{\alpha(t)} \quad ; \quad \Phi = \Gamma(1-\alpha(t))\ e^{-i\pi\alpha(t)}\ (\nu/\nu_o)^{\alpha(t)} \quad , \tag{4}$$

avec $\nu = (s-u)$.

Dans le cas d'une voie exotique, on a

$$(T - V) \;=\; 2\,g\,r \qquad (5.a)$$

et pour une voie non exotique :

$$(T + V) \;=\; 2\,g\,\Phi \quad . \qquad (5.b)$$

Le choix fait pour la conjugaison de charge du triplet $\vec{\pi}$ (Appendice A) et les équations (3.5) sont tout-à-fait cohérents ; en effet

$$\rho(\pi^{\pm} p) \;=\; \mp\,(F+D)\,(r-\Phi) \quad . \qquad (6)$$

Etant donné que

$$\sigma_{TOT}\,(\pi^- p) \;>\; \sigma_{TOT}\,(\pi^+ p) \quad ,$$

il faut que

$$\mathrm{Im}\;\rho_o(\pi^- p)\Big|_{t=0} \;>\; 0 \quad , \qquad (7)$$

où l'indice O est le bilan d'hélicité de la voie s , $n = |\lambda_3-\lambda_4-\lambda_1+\lambda_2|$. De (6) et (7) on conclut que $(F+D) < 0$. De la même façon de $\sigma_{TOT}(K^- p) > \sigma_{TOT}(K^+ p)$ il vient $\mathrm{Im}(A_2+\rho)_o\Big|_{t=0} > 0$ et donc $(F+D) < 0$.

Empiriquement il semble que les rapports $(F/D)_n = \alpha_n$ aient la même valeur pour les vecteurs et les tenseurs. Nous ferons l'hypothèse que ceci est vrai et nous pourrons obtenir les valeurs de F et D à partir des couplages du ρ et du f dans πN

$$\left.\begin{array}{l}\rho_o(\pi^- p) = (F+D)_o = -4.5 \\ f_o(\pi^- p) = (3F-D)_o = -19\end{array}\right\} \Longleftrightarrow \left\{\begin{array}{l}F_o = -5.875 \\ D_o = 1.375\end{array}\right\} \alpha_o = -4.3$$

(8)

$$\left.\begin{array}{l}\rho_1(\pi^- p) = (F+D)_1 = 17.9 \\ f_1(\pi^- p) = (3F-D)_1 = -1\end{array}\right\} \Longleftrightarrow \left\{\begin{array}{l}F_1 = 4.225 \\ D_1 = 13.675\end{array}\right\} \alpha_1 = 0.31 \quad .$$

Les valeurs obtenues pour α_n sont tout-à-fait comparables à celles indiquées dans les compilations de constantes de couplage[10]. On peut ainsi calculer les couplages du K_V et K_T dans les échanges d'hypercharge.

$$\begin{aligned} K_{V,T}(\pi N \to K\Sigma)_o &= -10.9 \quad ; \quad K_{V,T}(\pi N \to K\Sigma)_1 = -14.2 \\ K_{V,T}(\pi N \to K\Lambda)_o &= 8.1 \quad ; \quad K_{V,T}(\pi N \to K\Lambda)_1 = -13.2 \end{aligned} \qquad . \qquad (9)$$

On constate que le rapport $(\frac{n=0}{n=1})$ change de signe quand on passe du Σ au Λ comme baryon de l'état final, ce qui explique que les polarisations du Σ et du Λ sont sensiblement de signes opposés.

En outre, ce rapport est nettement plus grand que dans les échanges de charge, ce qui est clair sur les distributions angulaires à petit angle. La forme choisie pour un pôle de Regge est la suivante :

$$R_n(s,t) = g_n\left(-\frac{t'}{t_o}\right)^{n/2} \Gamma(1-\alpha(t))(1+\xi\, e^{-i\pi\alpha(t)})\,(\nu/\nu_o)^{\alpha(t)} \quad . \quad (10)$$

En utilisant les résultats des modèles duaux on peut écrire $(\nu/\nu_o)^{\alpha(t)} = (\alpha'_B\nu)^{\alpha(t)}\,\gamma^{\alpha(t)}$, où α'_B est la pente de la trajectoire des baryons de la voie s, que l'on prendra égale à $1\ \mathrm{GeV}^{-2}$. Les couplages g_n sont donnés par (9), il reste à déterminer la trajectoire $\alpha_k(t)$ et le paramètre d'échelle ν'_o. En tenant compte des masses et des largeurs du K_V et du K_T, on peut obtenir $\alpha_K(t)$ dans l'approximation linéaire

$$\alpha_K(t) = 0.37 + 0.79\ t \quad . \qquad (11)$$

L'hypothèse que SU(3) est exact à $t = 0$ nous conduit à poser :

$$\left\{ \gamma^{\alpha_M(o)} = \gamma'^{\alpha_K(o)} \right. \quad \text{avec par définition} \quad \left\{ \begin{aligned} \gamma &= \frac{\alpha'_B}{\nu_o} \\ \gamma' &= \frac{\alpha'_B}{\nu'_o} \end{aligned} \right. , \qquad (12)$$

où $\alpha_M(o)$, $\alpha_K(o)$ sont les ordonnées à l'origine des trajectoires (ρ, A_2)

et (K_V,K_T) et ν_o,ν'_o les paramètres d'échelle pour l'échange de (ρ,A_2) et de (K_V,K_T). Seul ν'_o est inconnu et peut donc être déterminé par (12)

$$\nu'_o = 0.175 \text{ GeV}^2 \quad . \tag{13}$$

Le fait que $\alpha_K(o) \neq \alpha_M(o)$ est une manifestation de la brisure de SU(3) par les masses. Dans (12), nous n'avons pas inclus les fonctions $\Gamma(1-\alpha(t))$ qui ont les pôles aux masses physiques qui brisent SU(3) . Les termes de pôles sont ainsi complètement définis sans nouveau paramètre.

B - Les termes d'absorption (PR et RP)

Les termes d'absorption sont obtenus par convolution des amplitudes d'échange du Pomeron et d'un pôle de Regge (convolution notée $*$). Le problème est ici de proposer une paramétrisation de l'échange du Pomeron. Nous avons choisi une forme très classique qui conduit à une section efficace totale constante, approximation valable pour les énergies qui nous intéressent. De plus cette amplitude ne contribue qu'à des termes correctifs , il ne serait pas réaliste de prendre une paramétrisation compliquée, incluant de nombreux paramètres. On masquerait ainsi totalement les effets des autres coupures. De même que dans le chapitre précédent, nous posons

$$P(A+B \to A+B) = i\, \sigma_\infty(AB)\, \nu\, e^{t[a(AB)+\alpha'_P \ell n \nu]}\, e^{-i\frac{\pi}{2}\alpha'_P t} \quad , \tag{14}$$

$\sigma_\infty(AB)$ est la valeur de la constante qui intervient dans une paramétrisation de la section efficace totale de la forme $\sigma_{TOT} = \sigma_\infty + G.s^{-1/2}$, le paramètre d'échelle $\nu_o = 1.\text{GeV}^2$ et la trajectoire du Pomeron est :

$$\alpha_P(t) = \alpha_P(o) + \alpha'_P t = 1. + 0.2t \quad . \tag{15}$$

Le Pomeron est supposé se découpler de l'hélicité flip. La pente du pic de diffraction est $2\,(a + \alpha'_P \ell n \nu)$, le terme $\alpha'_P \ell n \nu$ est une correction due à la pente du Pomeron qui donne le rétrécissement (skrinkage) du pic avec l'énergie. L'analogie avec la diffraction par un disque permet de relier $\sigma_\infty(AB)$ et $a(AB)$. Ainsi dans le cas d'un disque noir de

rayon R il vient :

$$\sigma_\infty = 2\pi R^2 \quad \text{et} \quad a = R^2/8 \ . \tag{16}$$

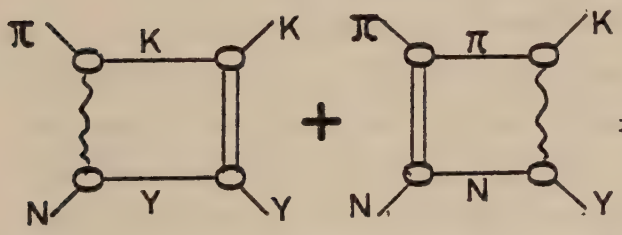
On peut introduire un facteur de transparence τ qui rend compte du fait que l'absorption n'est pas totale c'est-à-dire que l'on a ; $\sigma_\infty = 2\pi R^2(1-\tau)$. Si nous supposons que le rapport : $\sigma_\infty/a = 16\pi(1-\tau)$ est une constante, les paramètres ajustés dans πN nous donnent les valeurs de R et de τ. On trouve

$$\sigma_\infty(\pi N) = 21.5 \text{ mb} \qquad a = 2.75 \text{ GeV}^2$$

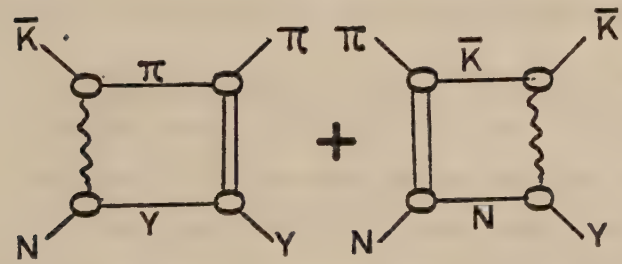
$$R = .927 \text{ fermi} \qquad \tau = .6 \ . \tag{17}$$

La valeur obtenue pour R est très proche du rayon hadronique conventionnel de 1 fermi.

Dans les réactions d'échange d'hypercharge suivant que l'on absorbe dans l'état initial où dans l'état final les Pomerons concernés sont ceux de réactions très différentes et peu faciles à mesurer. Pour $\pi N \to KY$ et $\bar{K}N \to \pi Y$ nous avons les contributions suivantes :



$$= (K_T+K_V)*P(KY)+P(\pi N) * (K_T+K_V) \tag{18}$$



$$= (K_T-K_V)*P(\pi Y) + P(\bar{K}N) * (K_T-K_V). \tag{19}$$

Afin de calculer ces différentes contributions, nous proposons un modèle naïf. Le Pomeron est composé d'une partie singlet de SU(3) $\tilde{P}$ qui se couple de la même façon à toutes les voies et d'une composante ayant les propriétés du f_o. $\sigma_\infty(\pi N)$ et $\sigma_\infty(KN)$ étant connus, toutes les $\sigma_\infty(AB)$

peuvent être calculées, de même que l'image géométrique et la donnée de $a(\pi N)$ déterminent tous les $a(AB)$ (Table I). Les termes d'absorption ne contiennent donc aucun paramètre libre, car nous nous plaçons dans le cadre des modèles d'absorption faible. Dans le chapitre précédent, la force de la coupure d'absorption était ajustée, mais le résultat de l'ajustement donnait pratiquement 1.

C - Les coupures Reggeon-Reggeon (R R)

Les règles de construction des diagrammes de dualité font que les coupures $R_1 R_2$ sont toujours du type $\Phi_1 * r_2$ ou $r_1 * \Phi_2$. Concrètement pour la réaction $\pi N \to KY$ tous les termes R.R seront de la forme $\Phi_s * r_M$ où l'indice s (resp. M) réfère à l'échange des mésons étranges K_V et K_T (resp. des mésons non étranges ρ, A_2, ω, f). Pour la voie exotique $\bar{K}N \to \pi Y$ nous aurons $r_s * \Phi_M$. Pour calculer les contributions des coupures R R. nous tenons compte de tous les états intermédiaires possibles. Dans les appendices A et B nous avons calculé tous les couplages nécessaires en fonction des quantités F et D déterminées ci-dessus. De nouveau, aucun paramètre n'est requis dans ce calcul. Nous devons tenir compte également des différents états d'hélicité de l'état intermédiaire (cf. équation A2). Si $n = |\lambda_3 - \lambda_4 - \lambda_1 + \lambda_2|$ est le bilan d'hélicité de la voie s et si $(\pm)$ dénote l'hélicité du baryon $(\pm 1/2)$, nous avons :

$$\begin{aligned} (M)_{n=0} &= (a_{++}) * (b_{++}) + (a_{+-}) * (b_{-+}) \\ (M)_{n=1} &= (a_{+-}) * (b_{--}) + (a_{++}) * (b_{+-}) \end{aligned} \quad , \qquad (20)$$

où a et b représentent les échanges des 2 Reggeons.

Tous les états intermédiaires pour les réactions considérées sont donnés en détail dans l'appendice C, ainsi que leur contributions respectives. En ce qui concerne les coupures comportant un Pomeron d'absorption dans l'état initial ou dans l'état final, le calcul est immédiat. Les diagrammes de quarks obéissent aux mêmes règles que celles des coupures R R. A nouveau, nous avons des Pomérons différents suivant la réaction et la place où l'on absorbe. Ainsi pour la réaction [a] $(\pi^+ p \to K^+ \Sigma^+)^\dagger$ nous avons

---

†Les notations sont celles de l'appendice C

$$P(\pi N) * [a_1] + [a_1] * P(K\Sigma) + P(\pi N) * [a_2] + [a_2] * P(K\Sigma) =$$

$$\{P(\pi N) + P(KN)\} * \{[a_1] + [a_2]\} = \{P(\pi N) + P(KN)\} * \{[a]\} \quad , \qquad (21)$$

où $\{[a]\}$ est la somme algébrique de toutes les contributions des coupures R.R. pour $\pi^+ p \to K^+\Sigma^+$. Avec les mêmes notations, nous pouvons écrire pour les autres réactions :

$$\bar{a} \ : \{P(KN) + P(\pi\Sigma)\} * \{[\bar{a}]\}$$

$$b \ : \{P(\pi N) + P(K\Sigma)\} * \{[b]\}$$

$$c \ : \{P(\pi N) + P(K\Sigma)\} * \{[c]\}$$

$$\bar{c} \ : \{P(KN) + P(\pi\Sigma)\} * \{[\bar{c}]\}$$

$$d \ : \{P(\pi N) + P(K\Lambda)\} * \{[d]\}$$

$$\bar{d} \ : \{P(KN) + P(\pi\Lambda)\} * \{[\bar{d}]\}$$

$$\bar{e} \ : \{P(KN) + P(\pi\Lambda)\} * \{[\bar{e}]\}$$

$$\bar{f} \ : \{P(KN) + P(\pi\Lambda)\} * \{[\bar{f}]\} \ .$$

D - Les coupures Reggeon-Pomeron-Reggeon (RPR)

Dans ces termes le Pomeron d'absorption intervient sur l'état intermédiaire, il y a donc lieu d'examiner toutes les contributions séparément. De plus, les règles de construction des diagrammes font que ces coupures sont toujours du type $\Phi_1 P \Phi_2$ ou $r_1 P r_2$ et jouent donc un rôle particulier (cf. le chapitre précédent). Nous donnons à titre d'exemple les contributions de ce type dans la réaction $[\bar{a}]\ K^- p \to \pi^- \Sigma^+$

$$RPR[\bar{a}] = \{P(KN)\} * \{[\bar{a}_1]\} + \{P(\pi\Sigma)\} * \{[\bar{a}_2] + [\bar{a}_4]\} + \{P(\pi\Lambda)\} * \{[\bar{a}_3]\} \ , \qquad (22)$$

où $[a_i]$ réfère aux notations de l'appendice C, les couplages sont inchangés mais les amplitudes de Regge sont de la forme $r_s r_M$ et non plus $r_s \Phi_M$. Pour les réactions du type $\pi N \to KY$ les termes sont du type $\Phi_s \Phi_M$. Il faut noter que les règles de sélection sur les diagrammes font que ces coupures RPR n'apparaissent pas si la voie t est exotique ce qui est le cas pour les réactions $[c]$ et $[\bar{c}]$, dans lesquelles contribuent exclusivement les termes R R et P R R.

E - Discussion des résultats obtenus

Les résultats sont présentés dans les figures 1 à 5. Il faut souligner que ce sont des prédictions, car nous n'avons aucun paramètre libre. Pour la réaction $\pi^+ p \to K^+\Sigma^+$, la structure en t de la section efficace différentielle est satisfaisante bien qu'à basse énergie, il semble y avoir un maximum secondaire que nous ne décrivons pas. Pour la dépendance en énergie, l'accord est moins bon pour $\pi^+ p \to K^+\Sigma^+$ (fig.1a) que pour les réactions $\pi^- p \to K^o\Sigma^o$ (fig.1b) et $\pi^- p \to K^o\Lambda$ (fig.1c). Ceci peut se comprendre aisément si l'on examine les trajectoires effectives. En paramétrisant les sections efficaces différentielles sous la forme

$$\frac{d\sigma}{dt} = f(t)\ S^{2\alpha_{eff}(t)-2} \quad , \qquad (23)$$

on peut extraire des résultats expérimentaux la trajectoire effective de l'échange de la voie t qui gouverne la dynamique de la réaction. Nous donnons dans le tableau ci-dessous les trajectoires effectives

Réaction	$\alpha(o)$	$\alpha'(o)$
[a] $\pi^+ p \to K^+\Sigma^+$ [10]	0.56	0.92
[b] $\pi^- p \to K^o\Sigma^o$ [11]	0.33 ± 0.06	0.23 ± 0.46
[d] $\pi^- p \to K^o\Lambda$ [11]	0.33 ± 0.06	0.43 ± 0.36

On constate que les réactions [b] et [d] ont une dépendance en énergie pour les petites valeurs de t très proche de celle qu'on attend des échanges K_V et K_T : $\alpha(o) \sim 0.33$, ce qui n'est pas du tout le cas de [a] . La présence de coupures d'absorption (même $\alpha(o)$ que le pôle) et de coupures R R ($\alpha(o) \sim 0$) ne modifie que peu $\alpha(o)$, par contre la pente de la trajectoire est affectée et tend à diminuer, les coupures ayant des $\alpha'(o)$ plus faibles que les pôles (typiquement $\alpha'_{coupure}(o) \sim 1/2\ \alpha'_{pôle}(o)$). Il n'est donc pas surprenant que [b] et [d] soient mieux décrites par le modèle que [a] . La polarisation dans les réactions $\pi^+ p \to K^+\Sigma^+$ (fig.3a) et $\pi^- p \to K^o\Lambda$ (fig.3b) est correctement reproduite. Pour les réactions renversées de lignes, à basse énergie l'accord obtenu est bon, ce qui indique que le mécanisme de brisure de la dégénérescence d'échange produit l'effet escompté. Pour $p_{lab} = 14.3$ GeV/c il y a

désaccord, la contribution des coupures disparaissant très vite quand l'énergie augmente. Néanmoins, il semble qu'il y ait un problème expérimental de normalisation relative entre $\pi^+p \to K^+\Sigma^+$[11] et $K^-p \to \pi^-\Sigma^+$[17] à cette énergie, car une étude avec le même appareillage de ces deux réactions[18] à $p_{lab} = 10$ GeV/c donne :

$$\left(\frac{d\sigma}{dt}\right)_{t=0} [\bar{a}] \;/\; \left(\frac{d\sigma}{dt}\right)_{t=0} [a] = 1.050 \pm 0.048 \quad , \qquad (24)$$

ce qui est en parfait accord avec les prédictions du modèle. En ce qui concerne les polarisations, les tendances sont reproduites (fig.3b), même si la description n'est pas parfaite. Il convient de souligner que de simples coupures d'absorption permettent de reproduire parfaitement les polarisations[19-20] et l'effet des singularités secondaires paraît plutôt néfaste à cet égard.

En ce qui concerne les réactions avec voie t exotique $\pi^-p \to K^+\Sigma^-[c]$ et $K^-p \to \pi^+\Sigma^-[\bar{c}]$ les résultats expérimentaux et les prédictions du modèle sont présentés dans les figures 4a.b. Il convient de faire quelques remarques sur les données expérimentales. Tout d'abord, ainsi que le suggère la philosophie des modèles d échange, les sections efficaces mesurées sont très petites[8]. Typiquement il y a 3 ordres de grandeur de différence entre les réactions avec échange de nombre quantique permis et celles avec échange "exotique". La variation du point de l'avant $\frac{d\sigma}{dt}(t=0)$ avec l'énergie pour $s \lesssim 7,8\ \text{GeV}^2$ est en s^{-10}, ce qui est difficilement explicable autrement que par des effets de voie s (ou de voie u, coupures BB par exemple). Pour des énergies supérieures le comportement change et devient compatible avec s^{-2} qui est ce que l'on attend de coupures Reggeon-Reggeon (R R). Outre cette dépendance en énergie, on observe que

$$\left[\frac{d\sigma}{dt}(K^-p \to \pi^+\Sigma^-) \;/\; \frac{d\sigma}{dt}(\pi^-p \to K^+\Sigma^-)\right]_{t=0} \gtrsim 30 - 50 \quad , \qquad (25)$$

pour $s \lesssim 7\text{-}8\ \text{GeV}^2$, ensuite les deux distributions angulaires sont sensiblement du même ordre de grandeur, de telles mesures pour $s \sim 10\text{-}15\ \text{GeV}^2$ seraient tres précieuses. Ces résultats indiquent que les forces de voie t n'apparaissent vraisemblablement qu'à partir de $p_{lab} \sim 3$ GeV/c . Il apparaît difficile de tester le calcul des coupures R.R à partir de ces réactions,

étant donnés le comportement surprenant et la pauvreté des résultats expérimentaux dans la région où les forces de voie t commencent à jouer. On peut néanmoins vérifier que l'ordre de grandeur donné par le modèle est correct et que la présence d'états intermédiaires différents pour les deux réactions explique qualitativement (25). On constate que le modèle donne un ordre de grandeur acceptable pour les deux réactions et rend compte assez bien de (25). Ces résultats confirment que l'approche utilisée a une certaine pertinence physique et que les termes de double échange interviennent dans la dynamique de ces réactions.

Les amplitudes de diffusion obtenues (Figs. 5a.b) sont très semblables à celles des analyses en amplitudes [2,3,7,21]. La partie imaginaire de l'échange vecteur dans le non flip présente un zéro périphérique très semblable à celui du ρ et du ω. La partie imaginaire s'annule avant la partie réelle et cette séparation des zéros vient bien des contributions des coupures secondaires. Nous avons montré au chapitre précédent que cet effet est suffisant pour décrire la polarisation $\pi^- p \to \pi^o n$. L'amplitude $(K_V)_{+-}$ et l'amplitude $(K_T)_{+-}$ ont un comportement de pôle de Regge ainsi que l'établissent les analyses en amplitudes. Dans la systématique que nous proposons dans la deuxième partie (chapitre VII) l'échange de tenseur à la forme suivante

$$T_{++} = \alpha(t) \times \text{Regge} .$$

C'est-à-dire que la partie imaginaire aura un zéro double à t_o tel que $\alpha(t_o) = 0$ et la partie réelle un zéro simple. Les amplitudes obtenues sont tout à fait compatibles avec un tel comportement.

En guise de conclusion on peut dire que le modèle proposé donne une description correcte des données expérimentales. Il est remarquable que sans paramètre libre on ait obtenu une normalisation absolue raisonnable pour les échanges d'exotiques. Un ajustement sur ces réactions constituerait une mesure de vertex de Gribov.

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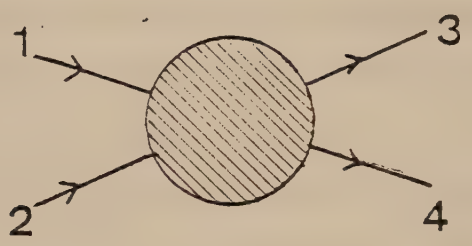
APPENDICE A

ROLE DE L'ISOSPIN

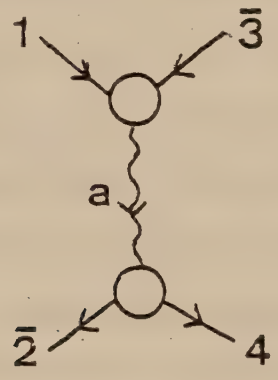
Dans cet appendice, nous donnons en détail les conventions choisies ainsi que les décompositions des amplitudes de diffusion en isospin de la voie t.

Echange d'un pôle de Regge

Considérons la réaction à 2 corps dont la voie s correspond à $1+2 \rightarrow 3+4$



Par définition l'échange d'un pôle de Regge (a) dans la voie t : $1+\bar{3} \rightarrow \bar{2}+4$ est représenté par le graphe suivant :



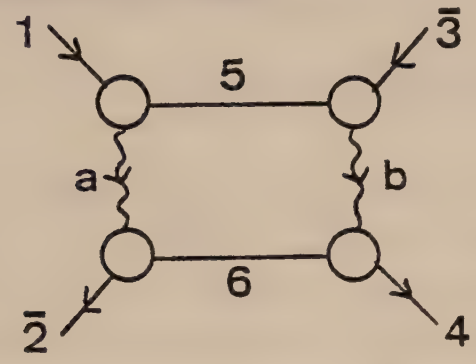
Une onde partielle de l'amplitude d'hélicité de voie s peut se décomposer en amplitudes d'isospin de la voie t : I.

$$M^J_{\lambda\mu}(s) = \langle \bar{2}4 | a^J_{\lambda\mu} | 1\bar{3} \rangle = \sum_I \langle \bar{2}4 | I \rangle \langle I | 1\bar{3} \rangle \langle I | a^J_{\lambda\mu} | I \rangle \quad , \qquad (A.1)$$

$\langle 2\bar{4} | I \rangle$ et $\langle I | 1\bar{3} \rangle$ sont les coefficients de Clebsh-Gordan correspondant à chacun des vertex, $|I\rangle$ dénote de façon simplifiée $|I,m\rangle$.

Le théorème de Wigner Eckart nous assure que $\langle I | a^J_{\lambda\mu} | I \rangle = [a^J_{\lambda\mu}]^I$, c'est-à-dire que cette amplitude ne dépend pas de m , projection de I . Cette amplitude est précisément l'amplitude d'onde partielle J entre des états d'hélicité en voie s λ et μ et d'isospin de voie t : I .

Echange d'une coupure de Regge



L'échange successif de deux pôles de Regge a et b se décompose de la façon suivante :

$$C^J_{\lambda\mu}(s) = \sum_{I_a, I_b, \sigma} \langle 1\bar{5} | I_a \rangle \langle 5\bar{3} | I_b \rangle \langle \bar{2}6 | I_a \rangle \langle \bar{6}4 | I_b \rangle \times [a_{\lambda\sigma}]^{I_a} [b_{\sigma\mu}]^{I_b} \quad (A.2),$$

σ est l'hélicité de l'état intermédiaire. Nous avons le produit des ondes partielles ce qui équivaut à faire une convolution sur les amplitudes complètes. Dans le cas d'une coupure d'absorption par exemple a ≡ P(Pomeron), b = R , il vient :

$$C^J_{\lambda\mu} = \sum_{I_b} \left(\langle 1\bar{3} | I_b \rangle \langle I_b | \bar{2}4 \rangle [b_{\sigma\mu}]^{I_b} \right) \cdot \left(\langle 1\bar{1} | 0 \rangle \langle 0 | \bar{2}2 \rangle P^J_{\lambda\sigma} \right) = \text{Pole}(I_b)^J \times$$

$$\times \text{(Pomeron)}^J .$$

Cette description en termes d'isospin de la voie t conduit à définir des états d'isospin pour les antiparticules, il y a quelques choix de phase à faire.

Conjugaison de charge et isospin

Soit un multiplet d'isospin $\vec{I}$ de dimension 2I+1 , sous l'action d'une rotation R de SU(2) les états $|I,m\rangle$ se transforment de la façon

suivante

$$R|I,m\rangle = \sum_{m'} \mathcal{D}^{I}_{m,m'}(R)|I,m'\rangle \quad , \tag{A.3}$$

où les $\mathcal{D}^{I}_{mm'}$ sont les matrices de rotation $(2I+1)\times(2I+1)$. Par définition le multiplet conjugué de charge obéit à la loi de transformation suivante

$$R|I,m\rangle_c = \sum_{m'} \mathcal{D}^{*I}_{mm'}(R)|I,m'\rangle_c \quad . \tag{A.4}$$

Or les coefficients de Clebsh-Gordan tabulés sont ceux de la réduction d'un produit de représentations associées à R et non a R^*. On peut néanmoins appliquer une transformation unitaire sur les états $|I,m\rangle_c$ pour définir de nouveaux états $|\overline{I,m}\rangle$ qui se transforment selon $\mathcal{D}^{I}_{mm'}(R)$ sous l'action d'une rotation R. Nous avons :

$$\mathcal{D}^{*}(R) = U^{-1}\mathcal{D}(R)U \quad , \qquad |\overline{I,m}\rangle = U|I,m\rangle_c$$

$$R|\overline{I,m}\rangle = \sum_{m'} \mathcal{D}^{I}_{mm'}(R)|\overline{I,m'}\rangle$$

la matrice U est telle que $U_{mm'} = \eta_I(-)^{I-m}\delta_{m-m}$ et donc

$$|\overline{I,m}\rangle = \eta_I(-)^{I-m}|I,-m\rangle_c \quad . \tag{A.5}$$

Le facteur de phase η_I est arbitraire, nous le choisirons de façon à rendre cohérente l'écriture des amplitudes. Par exemple si le Pomeron est un singlet de $SU(3)$, son couplage doit être le même dans toutes les réactions, ceci fixe certaines phases relatives à partir du choix fait pour des multiplets. Ainsi le triplet $\vec{\pi} = (\pi^+,\pi^o,\pi^-)$ est son propre C-conjugué , la désintégration $\pi^o \to \gamma+\gamma$ nous dit que $c = +1$ donc

$$C|1,0\rangle = |1,0\rangle \quad .$$

Si nous choisissons $\eta_I = +1$, il vient pour les états $|I,m\rangle$

$$|\overline{\pi_o}\rangle = -|\pi_o\rangle \qquad |\overline{\pi^{\pm}}\rangle = |\pi^{\mp}\rangle \tag{A.6}$$

De la même façon

$$|\bar{p}\rangle = |\bar{p}\rangle \ , \ |\bar{n}\rangle = -|\bar{n}\rangle \ ; \ |\overline{K^{\pm}}\rangle = |K^{\mp}\rangle \ , \ |\overline{K^{o}}\rangle = -|\overline{K^{o}}\rangle \ ; \ |\bar{\Lambda}\rangle = -|\bar{\Lambda}\rangle \ . \quad (A.7)$$

Le triplet $\vec{\Sigma}$ obéit à la même règle que $\vec{\pi}$.

Ces définitions nous permettent de calculer de façon inambigüe les décompositions en isospin des réactions qui nous intéressent.

$$\pi^{\pm} p \to \pi^{\pm} p = -\frac{1}{\sqrt{6}} I(0) \pm \frac{1}{2} I(1)$$

$$\pi^{-} p \to \pi^{o} n = \frac{1}{\sqrt{2}} I(1)$$

$$K^{\pm} p \to K^{\pm} p = \mp \frac{1}{2} I(0) + \frac{1}{2} I(1)$$

$$K^{\pm} n \to K^{\pm} n = \mp \frac{1}{2} I(0) - \frac{1}{2} I(1)$$

$$K^{o} p \to K^{o} p = -\frac{1}{2} I(0) - \frac{1}{2} I(1)$$

$$\overline{K^{o}} p \to \overline{K^{o}} p = \frac{1}{2} I(0) - \frac{1}{2} I(1)$$

$$K^{-} p \to \overline{K^{o}} n = I(1)$$

$$\pi^{-} \Sigma^{+} \to \pi^{-} \Sigma^{+} = \frac{1}{3} I(0) + \frac{1}{2} I(1)$$

$$\pi^{\pm} \Lambda \to \pi^{\pm} \Lambda = -\frac{1}{\sqrt{3}} I(0) = \pi^{o} \Lambda \to \pi^{o} \Lambda$$

$$\pi^{o} \Sigma^{o} \to \pi^{\pm} \Sigma^{\mp} = -\frac{1}{2} I(1)$$

$$\pi^{\pm} \Sigma^{o} \to \pi^{\pm} \Lambda = \mp \frac{1}{\sqrt{2}} I(1)$$

$$\pi^{o} \Sigma^{\pm} \to \pi^{\pm} \Lambda = \pm \frac{1}{\sqrt{2}} I(1) = \pi^{o} \Lambda \to \pi^{\mp} \Sigma^{\pm}$$

$$K^{+} \Sigma^{+} \to K^{+} \Sigma^{+} = \frac{1}{\sqrt{6}} I(0) - \frac{1}{2} I(1)$$

$$K^{\pm} \Sigma^{o} \to K^{\pm} \Sigma^{o} = \pm \frac{1}{\sqrt{6}} I(0)$$

$$K^{o} \Sigma^{o} \to K^{o} \Sigma^{o} = \frac{1}{\sqrt{6}} I(0)$$

$$\bar{K}^0 \Sigma^0 \rightarrow \bar{K}^0 \Sigma^0 = -\frac{1}{\sqrt{6}} I(0)$$

$$K^0 \Lambda \rightarrow K^0 \Lambda = -\frac{1}{\sqrt{2}} I(0)$$

$$\bar{K}^0 \Lambda \rightarrow \bar{K}^0 \Lambda = \frac{1}{\sqrt{2}} I(0)$$

$$K^0 \Sigma^0 \rightarrow K^+ \Sigma^- = -\frac{1}{\sqrt{2}} I(1)$$

$$K^0 \Sigma^0 \rightarrow K^0 \Lambda = \frac{1}{\sqrt{2}} I(1)$$

$$K^0 \Lambda \rightarrow K^+ \Sigma^- = -I(1)$$

$$\pi^+ p \rightarrow K^+ \Sigma^+ = -\frac{2}{3} I(1/2) + \frac{1}{3} I(3/2)$$

$$\pi^- p \rightarrow K^0 \Sigma^0 = -\frac{\sqrt{2}}{3} I(1/2) - \frac{\sqrt{2}}{3} I(3/2)$$

$$\pi^0 n \rightarrow K^0 \Sigma^0 = -\frac{1}{3} I(1/2) + \frac{2}{3} I(3/2)$$

$$\pi^0 n \rightarrow K^+ \Sigma^- = -\frac{\sqrt{2}}{3} I(1/2) + \frac{\sqrt{2}}{3} I(3/2)$$

$$\pi^- p \rightarrow K^0 \Lambda = \sqrt{\frac{2}{3}} I(1/2)$$

$$\pi^0 n \rightarrow K^0 \Lambda = -\frac{1}{\sqrt{3}} I(1/2)$$

$$K^- p \rightarrow \pi^0 \Sigma^0 = -\frac{1}{3} I(1/2) - \frac{2}{3} I(3/2)$$

$$K^- p \rightarrow \pi^- \Sigma^+ = \frac{2}{3} I(1/2) + \frac{1}{3} I(3/2)$$

$$\bar{K}^0 n \rightarrow \pi^- \Sigma^+ = I(3/2)$$

$$\bar{K}^0 n \rightarrow \pi^+ \Sigma^- = \frac{2}{3} I(1/2) + \frac{1}{3} I(3/2)$$

$$\bar{K}^0 p \rightarrow \pi^0 \Sigma^+ = -\frac{\sqrt{2}}{3} I(1/2) + \frac{\sqrt{2}}{3} I(3/2)$$

$$\bar{K}^0 p \rightarrow \pi^+ \Sigma^0 = \frac{\sqrt{2}}{3} I(1/2) - \frac{\sqrt{2}}{3} I(3/2)$$

$$\bar{K}^0 p \rightarrow \pi^+ \Lambda = -\sqrt{\frac{2}{3}} I(1/2)$$

$$K^- p \rightarrow \pi^0 \Lambda = \frac{1}{\sqrt{3}} I(1/2)$$

$$\bar{K}^{o} n \rightarrow \pi^{o} \Lambda = -\frac{1}{\sqrt{3}} I(1/2)$$

$$K^{-} n \rightarrow \pi^{-} \Lambda = \sqrt{\frac{2}{3}} I(1/2)$$

$$K^{-} n \rightarrow \pi^{-} \Sigma^{o} = \frac{\sqrt{2}}{3} I(1/2) - \frac{\sqrt{2}}{3} I(3/2)$$

$$K^{-} n \rightarrow \pi^{o} \Sigma^{-} = -\frac{\sqrt{2}}{3} I(1/2) + \frac{\sqrt{2}}{3} I(3/2).$$

APPENDICE B

SU(3) ET LES CONSTANTES DE COUPLAGES

Nous calculons les constantes de couplage des particules échangées dans les réactions considérées, en termes des facteurs isoscalaires de SU(3). Le vrai couplage s'obtient en précisant les états de charge (la charge de SU(2)) et en utilisant les décompositions données dans l'appendice A. En ce qui concerne les états $|8F\rangle$ et $|8D\rangle$ pour les couplages $M\bar{B}B$ nous convenons de les renormaliser respectivement par les facteurs $-\sqrt{6}$ et $-\frac{10}{\sqrt{30}}$, par définition nous poserons également $\rho\bar{K}K = +1$ et $A_2\bar{K}K = -1$. Nous obtenons ainsi les expressions simplifiées par exemple :

$$\rho(\pi^{\pm} p \to \pi^{\pm} p) = \mp (F + D) \quad ,$$

en tenant compte du facteur d'isospin $\pm \frac{1}{2}$.

a) $\underline{\pi N \to \pi N}$

$$\rho\pi\pi = -2 \qquad \rho\bar{N}N = (F + D) \qquad \rho(\pi N) = -2(F + D) \; .$$

Le $f_o(J^P = 2^+)$ est en fait un mélange de 2 états purs de SU(3) notés f_1 et f_8. Si l'on fait l'hypothèse que le f'_o (1514) est un état pur $(S\bar{S})$ on obtient les angles "magiques" du modèle des quarks, le mélange est alors :

$$f_o = \frac{1}{\sqrt{3}} \left[\sqrt{2}\, f_1 + f_8 \right] \qquad f'_o = \frac{1}{\sqrt{3}} \left[\sqrt{2}\, f_8 - f_1 \right] \; .$$

Les couplages du f_1 sont inconnus, on peut les relier à ceux du f_8 en imposant certaines contraintes physiques, telle que la règle de Zweig exacte, c'est-à-dire

$$f'_o\pi\pi = 0 \qquad \text{et} \qquad f'_o \bar{N}N = 0 \; .$$

Nous avons alors :

$$f_o\pi\pi = \sqrt{6} \qquad f_o\bar{N}N = D-3F \qquad f_o(\pi N) = \sqrt{6}\,(D-3F) \; .$$

b) $\underline{KN \to KN \quad , \quad \bar{K}N \to \bar{K}N}$

$$\rho K\bar{K} = \omega K\bar{K} = -1 \qquad f_o K\bar{K} = A_2 K\bar{K} = +1$$

$$\rho(KN) = -(F+D) \qquad A_2(KN) = (F+D)$$

$$\rho(\bar{K}N) = (F+D) \qquad A_2(\bar{K}N) = (F+D) \ .$$

Pour le ω , on a encore le mélange (ω,φ) , en imposant $\varphi N\bar{N} = 0$ et $\varphi_1 MM = 0$ car le couplage n'est pas invariant par conjugaison de charge, il vient :

$$f_o(KN) = (D-3F) \qquad \omega(KN) = -(D-3F)$$

$$f_o(\bar{K}N) = -(D-3F) \qquad \omega(\bar{K}N) = -(D-3F) \ .$$

c) $\underline{\pi\Sigma \to \pi\Sigma \ , \ K\Sigma \to K\Sigma \ , \ \bar{K}\Sigma \to \bar{K}\Sigma}$

$$f_o \bar{\Sigma}\Sigma = \sqrt{6}F \qquad \rho\, \bar{\Sigma}\Sigma = -2F$$

$$f_o(\pi\Sigma) = 6F \qquad \rho(\pi\Sigma) = 4F$$

$$f_o(K\Sigma) = -\omega(K\Sigma) = \sqrt{6}F \qquad \rho(K\Sigma) = -A_2(K\Sigma) = 2F$$

$$f_o(\bar{K}\Sigma) = \omega(\bar{K}\Sigma) = -\sqrt{6}F \qquad \rho(\bar{K}\Sigma) = A_2(\bar{K}\Sigma) = -2F \ .$$

d) $\underline{\pi\Lambda \to \pi\Lambda \quad , \quad K\Lambda \to K\Lambda \quad , \quad \bar{K}\Lambda \to \bar{K}\Lambda}$

$$f_o \bar{\Lambda}\Lambda = \frac{\sqrt{2}}{3}[2D-3F] = \omega\bar{\Lambda}\Lambda \qquad f_o(\pi\Lambda) = \frac{2}{\sqrt{3}}(2D-3F) \ ,$$

$$f_o(K\Lambda) = \frac{\sqrt{2}}{3}(2D-3F) \qquad \omega(K\Lambda) = -\frac{\sqrt{2}}{3}(2D-3F)$$

$$f_o(\bar{K}\Lambda) = -\frac{\sqrt{2}}{3}(2D-3F) \qquad \omega(\bar{K}\Lambda) = -\frac{\sqrt{2}}{3}(2D-3F)$$

e) $\underline{\pi\Lambda \to \pi\Sigma \quad , \quad K\Lambda \to K\Sigma \quad , \quad \bar{K}\Lambda \to K\Sigma}$

$$\rho\Lambda\Sigma = -\sqrt{\frac{2}{3}}D \qquad \rho(\pi\Lambda \to \pi\Sigma) = 2\sqrt{\frac{2}{3}}D$$

$$\rho(K\Lambda \to K\Sigma) = \sqrt{\frac{2}{3}}D \qquad A_2(K\Lambda \to K\Sigma) = -\sqrt{\frac{2}{3}}D$$

$$\rho(\bar{K}\Lambda \to \bar{K}\Sigma) = -\sqrt{\frac{2}{3}}D \qquad A_2(\bar{K}\Lambda \to \bar{K}\Sigma) = -\sqrt{\frac{2}{3}}D$$

f) $\underline{\pi N \to K\Sigma/\Lambda \quad ; \quad \bar{K}N \to \pi\Sigma/\Lambda}$

$$\bar{K}_T \pi \bar{K} = -\frac{3}{\sqrt{6}} \qquad \bar{K}_T \bar{N}\Sigma = \frac{3}{\sqrt{6}}\,(D-F) \qquad \bar{K}_T \bar{N}\Lambda = \frac{1}{\sqrt{6}}\,(D+3F)$$

$$K_V(\pi N \to K\Sigma) = \frac{3}{2}\,(F-D) = K_T(\pi N \to K\Sigma)$$

$$K_V(\pi N \to K\Lambda) = -\frac{1}{2}\,(D+3F) = K_T(\pi N \to K\Lambda)$$

$$K_V(\bar{K}N \to \pi\Sigma) = \frac{3}{2}\,(F-D) = -K_T(\bar{K}N \to \pi\Sigma)$$

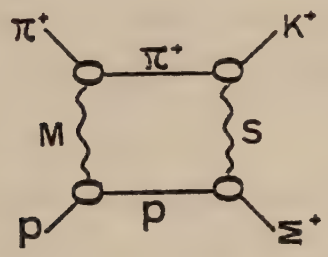
$$K_V(\bar{K}N \to \pi\Lambda) = -\frac{1}{2}\,(D+3F) = -K_T(\bar{K}N \to \pi\Lambda) \quad .$$

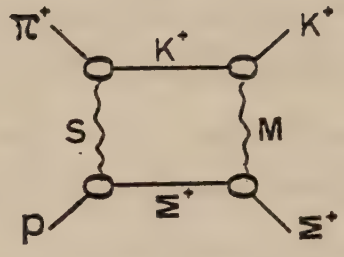
APPENDICE C

Dans cet appendice, nous précisons les états intermédiaires que nous introduisons dans le calcul des coupures R R. Il faut noter que nous négligerons dans ces états intermédiaires les contributions du η et du Δ qui sont peu couplés. Pour les calculs nous utilisons les résultats des appendices A et B . Dans le premier cas étudié $\pi^+p \to K^+\Sigma^+$ nous précisons les notations auxquelles nous nous tiendrons par la suite

$$a)\quad \pi^+p \;\to\; K^+\Sigma^+ \quad .$$

Il n'y a ici que deux possibilités, les K_V et K_T sont échangés en premier ou en second lieu





$[a_1] = (\pi^+p \to \pi^+p)*(\pi^+p \to K^+\Sigma^+)$ $\quad [a_2] = (\pi^+p \to K^+\Sigma^+)*(K^+\Sigma^+ \to K^+\Sigma^+)$

Le calcul des couplages nous donne :

	$\pi^+p \to K^+\Sigma^+$	$S_{(++)}{}^M{}_{(++)}$	$S_{(+-)}{}^M{}_{(-+)}$	$S_{(++)}{}^M{}_{(+-)}$	$S_{(+-)}{}^M{}_{(--)}$
a_1	$(\pi^+p \to \pi^+p)*(\pi^+p \to K^+\Sigma^+)$ $-4(F-D)\,\Phi_s(F-D)\,r_M$	-210.3	357.2	-274.	-274.
a_2	$(\pi^+p \to K^+\Sigma^+)*(K^+\Sigma^+ \to K^+\Sigma^+)$ $-8(F-D)\,\Phi_s\,F\,r_M$	-340.8	-319.2	245.	-444.

la notation SM réfère au fait que les deux Reggeons échangés sont très différents l'un est un méson non étrange M (ρ, A_2, f, ω), l'autre transporte de l'hypercharge S(K_V, K_T) .

Numériquement, nous aurons quatre quantités à calculer pour chaque coupure, concrètement $S_{(++)}M_{(++)}$ et $S_{(+-)}M_{(-+)}$ pour l'amplitude ayant $n = 0$, $S_{(++)}M_{(+-)}$ et $S_{(+-)}M_{(--)}$ pour la composante $n = 1$. Nous noterons les réactions $\bar{K}N \to \pi Y$ par une lettre surmontée d'une barre par référence à un renversement de lignes pour les obtenir à partir de la voie $\pi N \to KY$. Les valeurs numériques tabulées sont les valeurs des couplages

	$K^-p \to \pi^-\Sigma^+$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
$\bar{a}1$	$(K^-p \to K^-p)*(K^-p \to \pi^-\Sigma^+)$ $-8(F-D)r_s F \Phi_M$	-340.8	-319.4	245.	-444.
$\bar{a}2$	$(K^-p \to \pi^o\Sigma^o)*(\pi^o\Sigma^o \to \pi^-\Sigma^+)$ $-2(F-D)r_s F \Phi_M$	- 85.2	- 79.9	61.3	-111.
$\bar{a}3$	$(K^-p \to \pi^o\Lambda)*(\pi^o\Lambda \to \pi^-\Sigma^+)$ $\frac{2}{3}(D+3F)r_s D \Phi_M$	-14.9	-240.4	-148.2	24.2
$\bar{a}4$	$(K^-p \to \pi^-\Sigma^+)*(\pi^-\Sigma^+ \to \pi^-\Sigma^+)$ $-8(F-D)r_s F \Phi_M$	-340.8	-319.4	245.	-444.

	$\pi^-p \to K^o\Sigma^o$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
b1	$(\pi^-p \to \pi^-p)*(\pi^-p \to K^o\Sigma^o)$ $-4\sqrt{2}(F-D)\Phi_s F r_M$	-2 41.	-225.9	173.3	-314.
b2	$(\pi^-p \to K^o\Sigma^o)*(K^o\Sigma^o \to K^o\Sigma^o)$ $-2\sqrt{2}(F-D)\Phi_s F r_M$	-120.5	-113.	86.7	-157.
b3	$(\pi^-p \to K^o\Lambda)*(K^o\Lambda \to K^o\Sigma^o)$ $\frac{2}{3}\sqrt{2}(D+3F)\Phi_s D r_M$	-20.9	-336.5	-207.5	33.8
b4	$(\pi^-p \to \pi^o n)*(\pi^o n \to K^o\Sigma^o)$ $\sqrt{2}(F-D)\Phi_s(F+D) r_M$	46.1	239.2	-183.5	60.1

	$\pi^- p \to K^+\Sigma^-$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
c1	$(\pi^- p \to \pi^o n) * (\pi^o n \to K^+\Sigma^-)$ $-2(F-D)\Phi_s (F+D) r_M$	65.3	338.3	-259.5	85.
c2	$(\pi^- p \to K^o\Sigma^o) * (K^o\Sigma^o \to K^+\Sigma^-)$ $-4(F-D)\Phi_s \; F \; r_M$	-170.4	-159.7	122.5	-222.1
c3	$(\pi^- p \to K^o\Lambda) * (K^o\Lambda \to K^+\Sigma^-)$ $-\frac{4}{3}(D+3F)\Phi_s \; Dr_M$	29.8	480.7	296.4	-48.3

	$K^- p \to \pi^-\Sigma^-$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
$\bar{c}1$	$(K^- p \to \bar{K}^o n) * (\bar{K}^o n \to \pi^+\Sigma^-)$ $-4(F-D) r_s (F+D) \Phi_M$	-130.5	-676.6	519.	-170.
$\bar{c}2$	$(K^- p \to \pi^o\Sigma^o) * (\pi^o\Sigma^o \to \pi^+\Sigma^-)$ $-2(F-D) r_s \; F \; \Phi_M$	-85.2	-80.	61.3	-111.
$\bar{c}3$	$(K^- p \to \pi^o\Lambda) * (\pi^o\Lambda \to \pi^+\Sigma^-)$ $-\frac{2}{3}(D+3F) r_s \; D \; \Phi_M$	14.9	+240.4	148.2	-24.1

	$\pi^- p \rightarrow K^o \Lambda$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
d1	$(\pi^- p \rightarrow \pi^- p) * (\pi^- p \rightarrow K^o \Lambda)$ $-4\sqrt{\frac{2}{3}}\,(D+3F)\Phi_s . F\, r_M$	-311.9	363.7	224.3	505.7
d2	$(\pi^- p \rightarrow \pi^o n) * (\pi^o n \rightarrow K^o \Lambda)$ $-\sqrt{\frac{2}{3}}\,(D+3F)\Phi_s (F+D) r_M$	-59.7	385.1	237.5	96.8
d3	$(\pi^- p \rightarrow K^o \Lambda) * (K^o \Lambda \rightarrow K^o \Lambda)$ $\frac{2}{3}\sqrt{\frac{2}{3}}\,(D+3F)\Phi_s (2D-3F) r_M$	-180.3	-210.5	-129.8	292.3
d4	$(\pi^- p \rightarrow K^o \Sigma^o) * (K^o \Sigma^o \rightarrow K^o \Lambda)$ $2\sqrt{\frac{2}{3}}\,(F-D)\,\Phi_s\, D\, r_M$	-16.3	211.	-161.9	-21.2

	$\bar{K}^o p \rightarrow \pi^+ \Lambda$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
$\bar{d}1$	$(\bar{K}^o p \rightarrow \bar{K}^o p) * (\bar{K}^o p \rightarrow \pi^+ \Lambda)$ $-2\sqrt{\frac{2}{3}}\,(D+3F) r_s (F-D) \Phi_M$	-192.4	-406.6	-250.8	312.
$\bar{d}2$	$(\bar{K}^o p \rightarrow \pi^o \Sigma^+) * (\pi^o \Sigma^+ \rightarrow \pi^+ \Lambda)$ $2\sqrt{\frac{2}{3}}\,(F-D)\, r_s\, D\, \Phi_M$	-16.3	211.	-161.9	-21.2
$\bar{d}3$	$(\bar{K}^o p \rightarrow \pi^+ \Lambda) * (\pi^+ \Lambda \rightarrow \pi^+ \Lambda)$ $\frac{2}{3}\sqrt{\frac{2}{3}}\,(D+3F) r_s (2D-3F) \Phi_M$	-180.3	-210.5	-129.8	292.3
$\bar{d}4$	$(\bar{K}^o p \rightarrow \pi^+ \Sigma^o) * (\pi^+ \Sigma^o \rightarrow \pi^+ \Lambda)$ $2\sqrt{\frac{2}{3}}\,(F-D) r_s\, D\, \Phi_M$	-16.3	211.	-161.9	-21.2

	$K^- p \to \pi^0 \Lambda$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
$\bar{e}1$	$(K^- p \to K^- p) * (K^- p \to \pi^0 \Lambda)$ $\frac{4}{\sqrt{3}} (D+3F) r_s \; F \; \Phi_M$	220.6	-257.2	-158.6	-357.6
$\bar{e}2$	$(K^- p \to \bar{K}^0 n) * (\bar{K}^0 n \to \pi^0 \Lambda)$ $-\frac{2}{\sqrt{3}} (D+3F) r_s \; (F+D) \Phi_M$	-84.4	544.6	335.9	136.9
$\bar{e}3$	$(K^- p \to \pi^0 \Lambda) * (\pi^0 \Lambda \to \pi^0 \Lambda)$ $-\frac{2}{3\sqrt{3}} (D+3F) r_s (F+D) \; \Phi_M$	127.5	148.5	92.	-206.7
$\bar{e}4$	$(K^- p \to \pi^- \Sigma) * (\pi^- \Sigma^+ \to \pi^0 \Lambda)$ $-\frac{4}{\sqrt{3}} (F-D) \; r_s \; D \; \Phi_M$	23.	-298.4	229.	30.

	$K^- n \to \pi^- \Lambda$	$S_{(++)}M_{(++)}$	$S_{(+-)}M_{(-+)}$	$S_{(++)}M_{(+-)}$	$S_{(+-)}M_{(--)}$
$\bar{f}1$	$(K^- n \to K^- n) * (K^- n \to \pi^- \Lambda)$ $2\sqrt{\frac{2}{3}} (D+3F) r_s (F-D) \Phi_M$	192.4	406.6	250.8	-312.
$\bar{f}2$	$(K^- n \to \pi^- \Lambda) * (\pi^- \Lambda \to \pi^- \Lambda)$ $-\frac{2}{3}\sqrt{\frac{2}{3}} (D+3F) r_s (2D-3F) \Phi_M$	180.3	210.5	129.8	-292.3
$\bar{f}3$	$(K^- n \to \pi^- \Sigma^0) * (\pi^- \Sigma^0 \to \pi^- \Lambda)$ $-2\sqrt{\frac{2}{3}} (F-D) r_s \; D \; \Phi_M$	16.3	-211.	161.9	21.2
$\bar{f}4$	$(K^- n \to \pi^0 \Sigma^-) * (\pi^0 \Sigma^- \to \pi^- \Lambda)$ $-2\sqrt{\frac{2}{3}} (F-D) r_s \; D \; \Phi_M$	16.3	-211.	161.9	21.2

TABLE I

$\sigma_\infty(\pi N) = \tilde{P} + (3F-D) = 21.5^*$ mb	$a(\pi N) = 2.75^*$ GeV^{-2}
$\sigma_\infty(KN) = \tilde{P} + \left(\frac{3F-D}{2}\right) = 18^*$ mb	$a(KN) = 2.30$ GeV^{-2}
$\sigma_\infty(\pi\Sigma) = \tilde{P} + 2F = 18.5$ mb	$a(\pi\Sigma) = 2.37$ GeV^{-2}
$\sigma_\infty(\pi\Lambda) = \tilde{P} - \frac{2}{3}(2D-3F) = 19.8$ mb	$a(\pi\Lambda) = 2.54$ GeV^{-2}
$\sigma_\infty(K\Sigma) = \tilde{P} + F = 16.5$ mb	$a(K\Sigma) = 2.11$ GeV^{-2}
$\sigma_\infty(K\Lambda) = \tilde{P} - \frac{(2D-3F)}{2} = 17.2$ mb	$a(K\Lambda) = 2.20$ GeV^{-2}
$\sigma_\infty(NN) = \tilde{P} + \frac{(3F-D)^2}{2} = 38$ mb	$a(NN) = 4.86$ GeV^{-2}
$\sigma_\infty(N\Sigma) = \tilde{P} + F(D-3F) = 28.5$ mb	$a(N\Sigma) = 3.63$ GeV^{-2}
$\sigma_\infty(N\Lambda) = \tilde{P} + \frac{1}{3}(D-3F)(2D-3F) = 33.2$mb	$a(N\Lambda) = 4.23$ GeV^{-2} .

Pour les hadrons, nous avons supposé $\tau = .6$, constant. Les quantités munies d'un astérisque sont les valeurs prises comme références. On peut constater que la prédiction pour $\sigma_\infty(NN)$ est satisfaisante de même que pour σ_∞ (NY) où les valeurs expérimentales donnent 30-33 mb.

fig 1-a

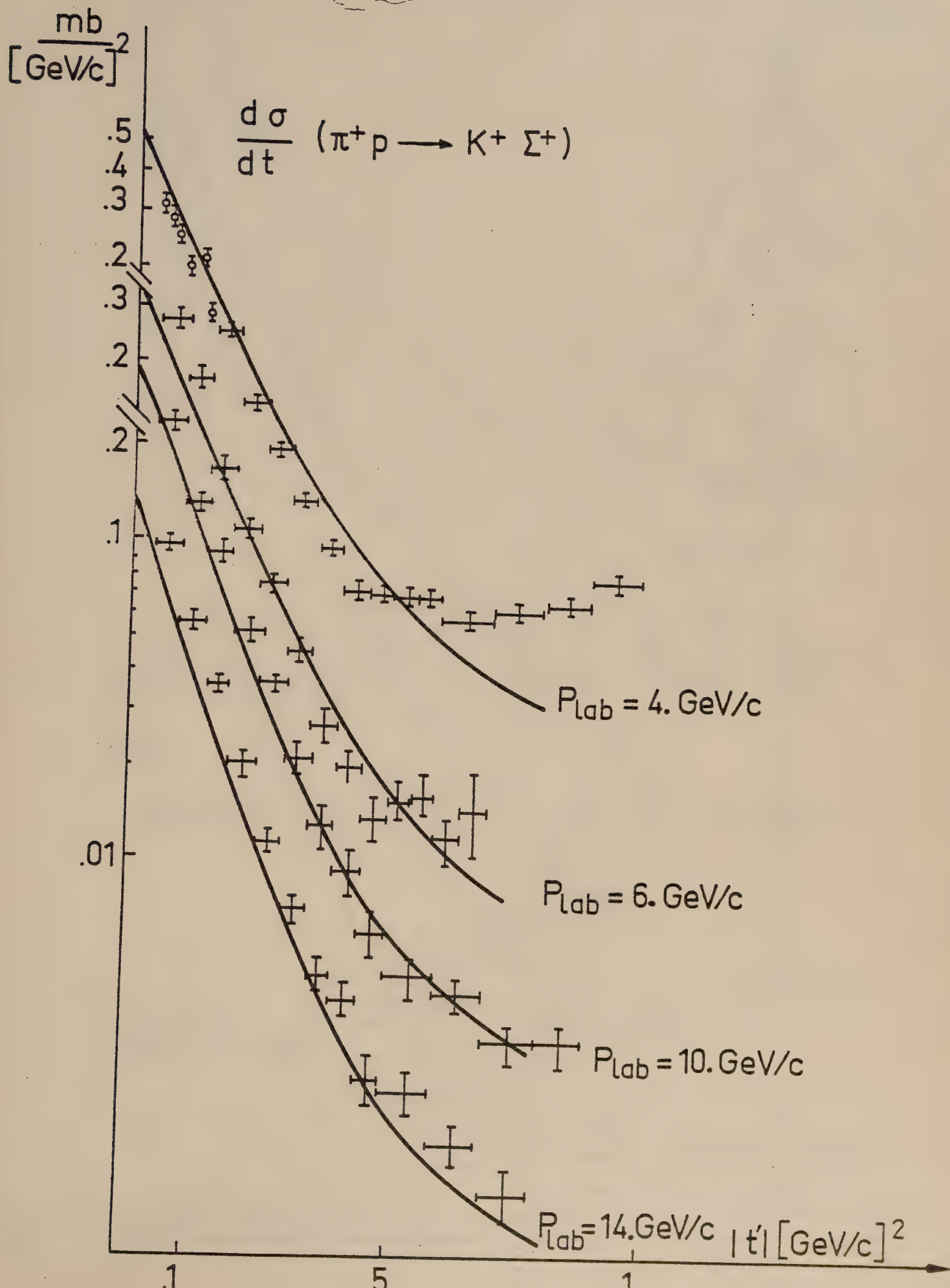
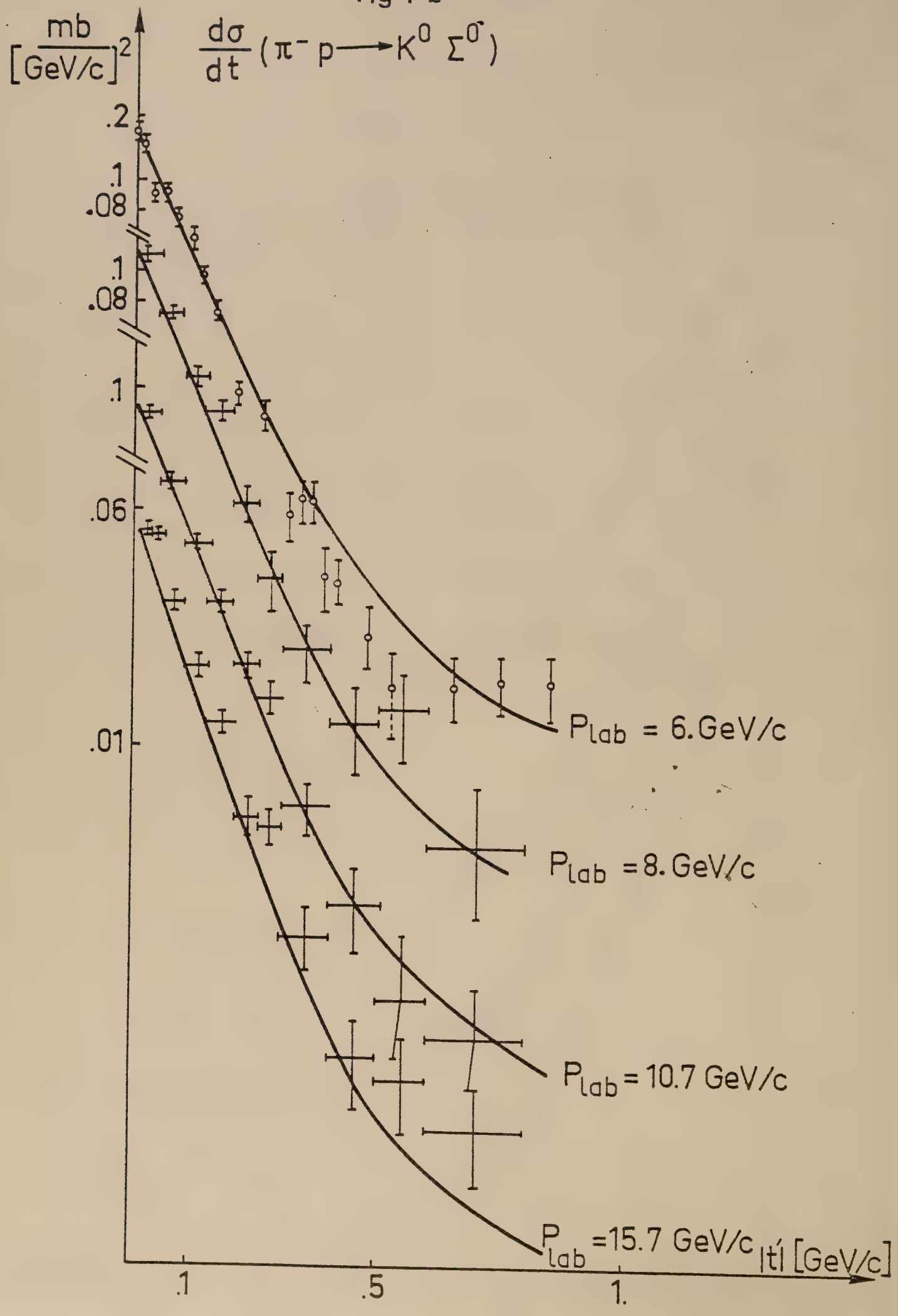

mb / [GeV/c]^2
dσ/dt (π+p → K+ Σ+)
.5
.4
.3
.2
.3
.2
.2
.1
.01
Plab = 4. GeV/c
Plab = 6. GeV/c
Plab = 10. GeV/c
Plab = 14. GeV/c
|t'| [GeV/c]^2
.1
.5
1.

fig 1-b


mb / [GeV/c]²
dσ/dt (π⁻ p → K⁰ Σ⁰)
.2
.1
.08
.1
.08
.1
.06
.01
P_lab = 6. GeV/c
P_lab = 8. GeV/c
P_lab = 10.7 GeV/c
P_lab = 15.7 GeV/c
.1
.5
1.
|t| [GeV/c]²

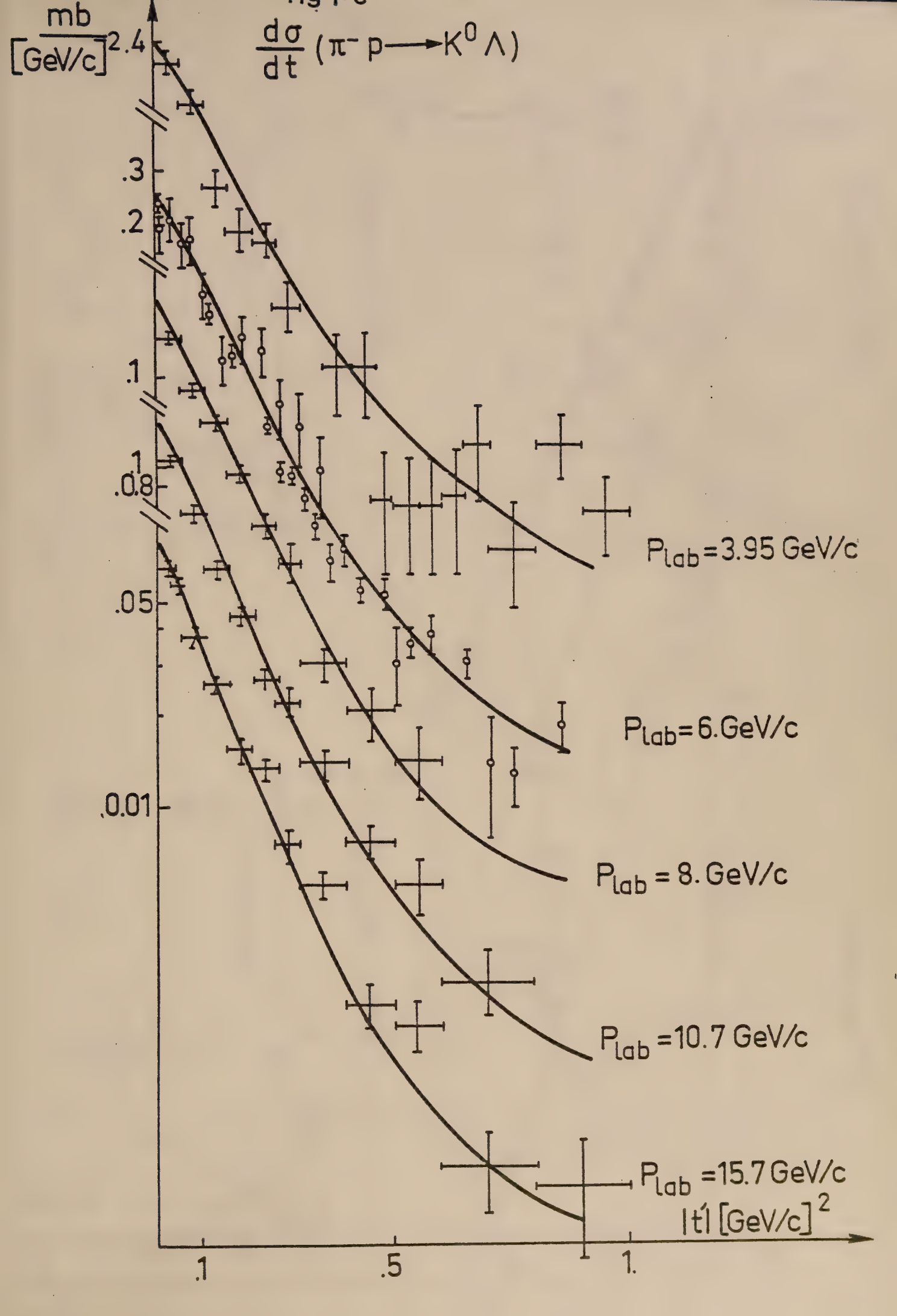
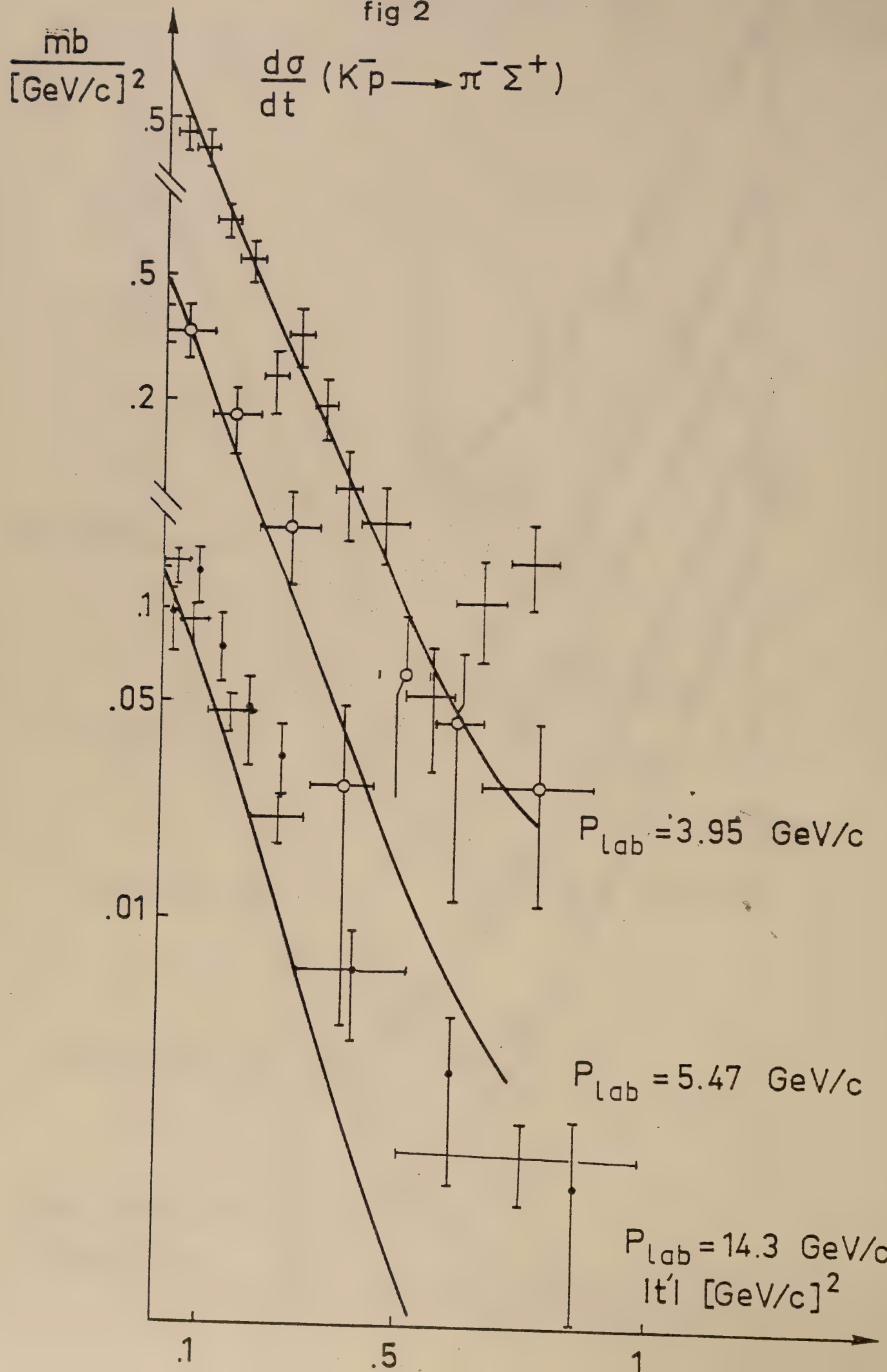
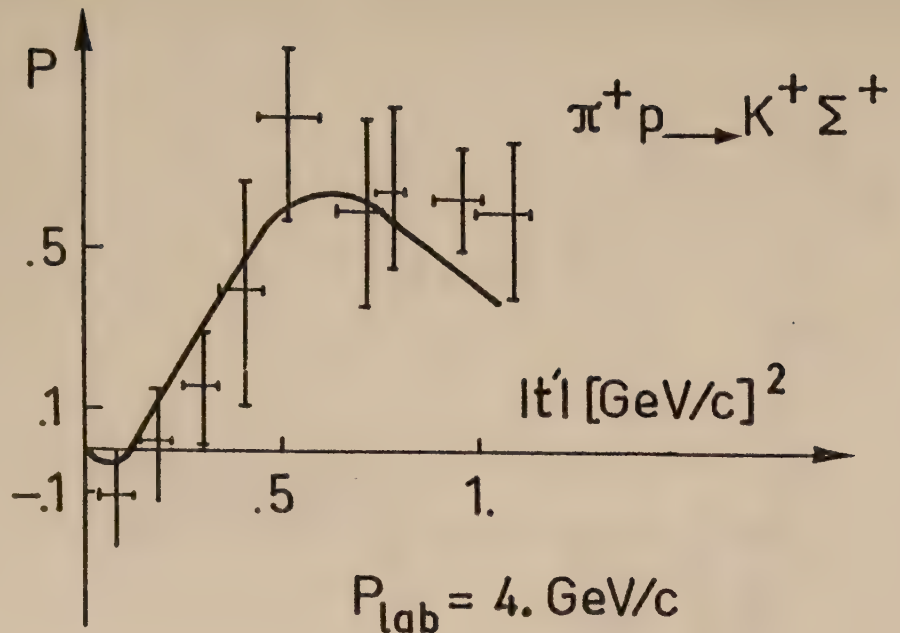
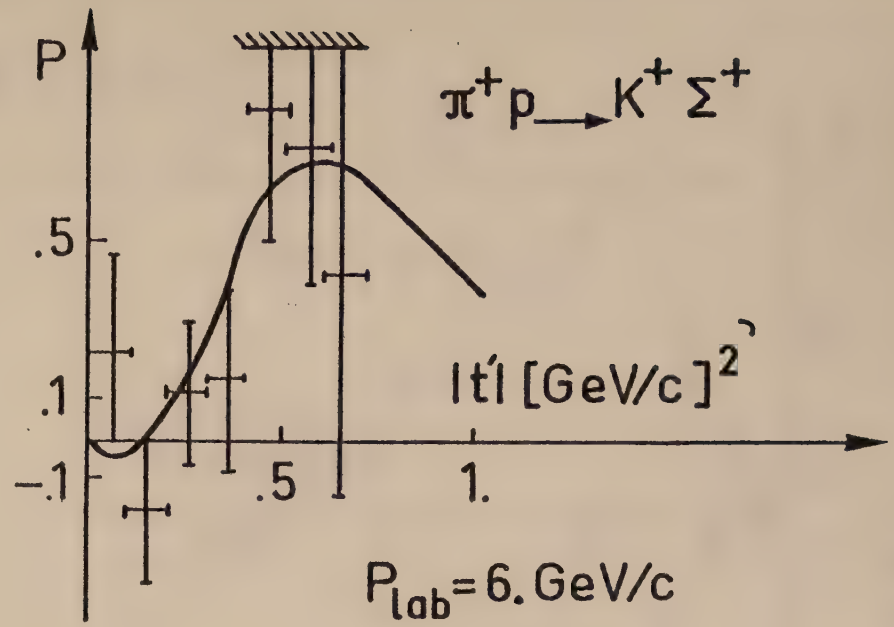

mb
[GeV/c]²
dσ/dt (π⁻ p → K⁰ Λ)
2.4
.3
.2
.1
.08
.05
.01
P_lab = 3.95 GeV/c
P_lab = 6. GeV/c
P_lab = 8. GeV/c
P_lab = 10.7 GeV/c
P_lab = 15.7 GeV/c
|t| [GeV/c]²
.1
.5
1.

fig 2


mb / [GeV/c]²
dσ/dt (K⁻p → π⁻Σ⁺)
.5
.5
.2
.1
.05
.01
P_lab = 3.95 GeV/c
P_lab = 5.47 GeV/c
P_lab = 14.3 GeV/c
|t'| [GeV/c]²
.1
.5
1

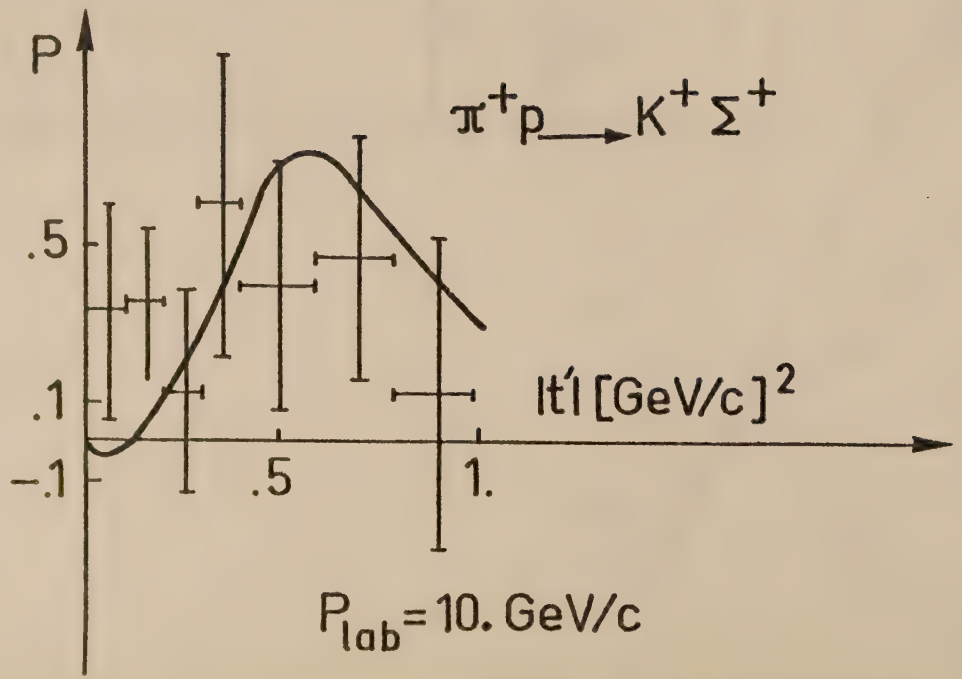


P_{lab} = 4. GeV/c

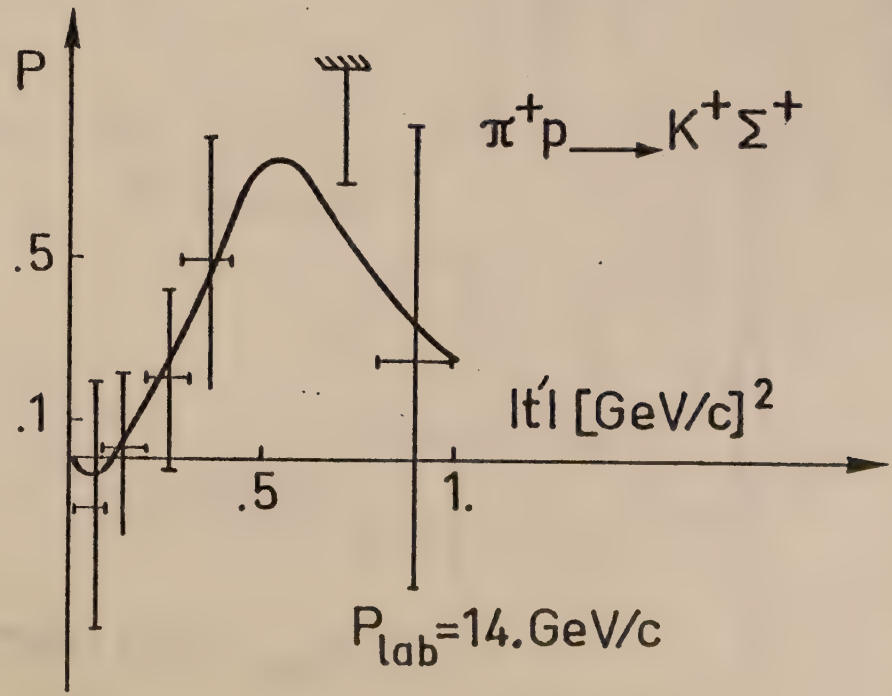


P_{lab} = 6. GeV/c

fig 3-a



P_{lab} = 10. GeV/c



P_{lab} = 14. GeV/c

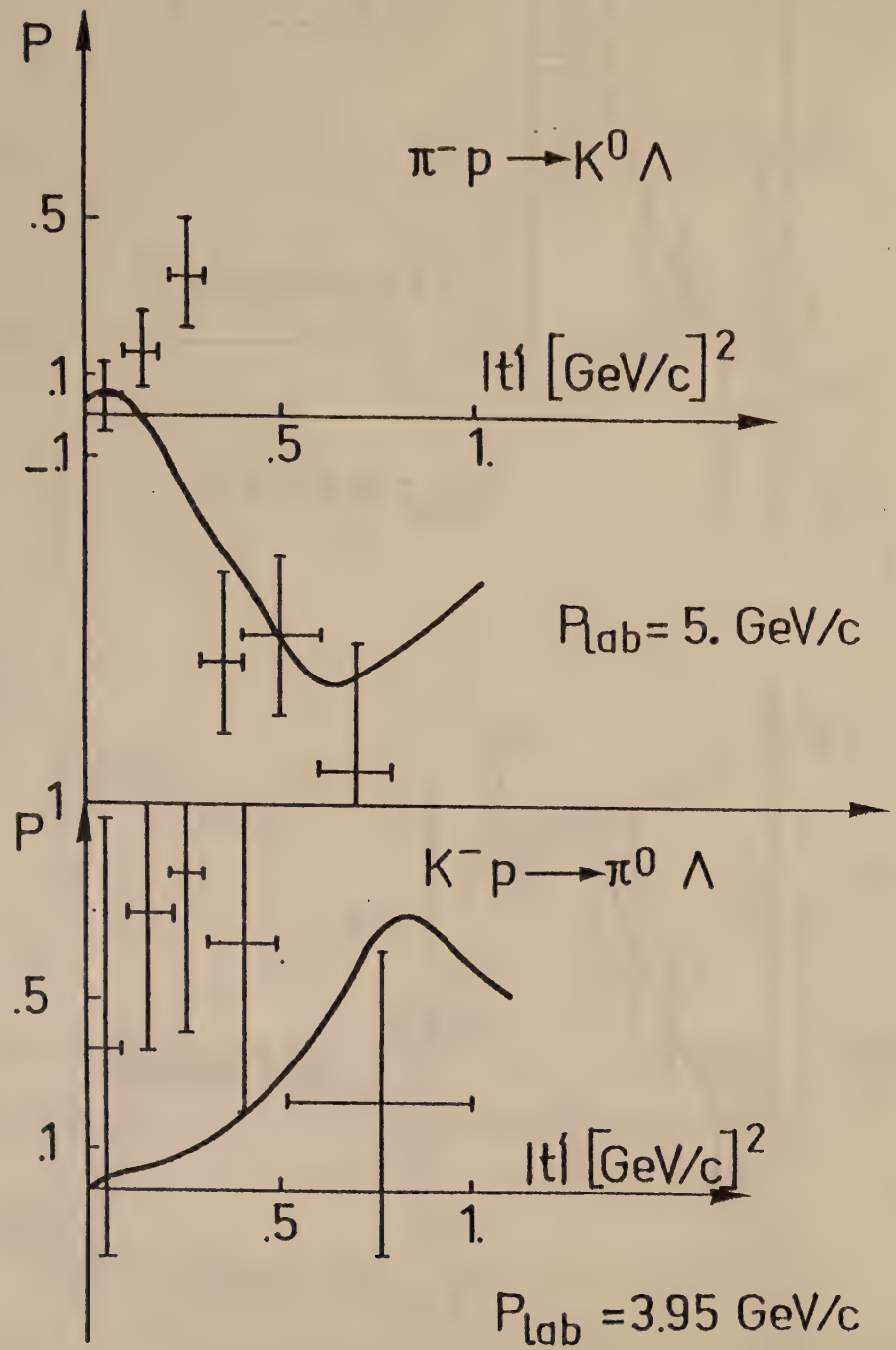

P
$\pi^- p \rightarrow K^0 \Lambda$
.5
.1
-.1
$|t'|$ [GeV/c]^2
.5
1.
P_{lab}= 5. GeV/c
P^1
$K^- p \rightarrow \pi^0 \Lambda$
.5
.1
$|t'|$ [GeV/c]^2
.5
1.
P_{lab} =3.95 GeV/c

fig 3-b

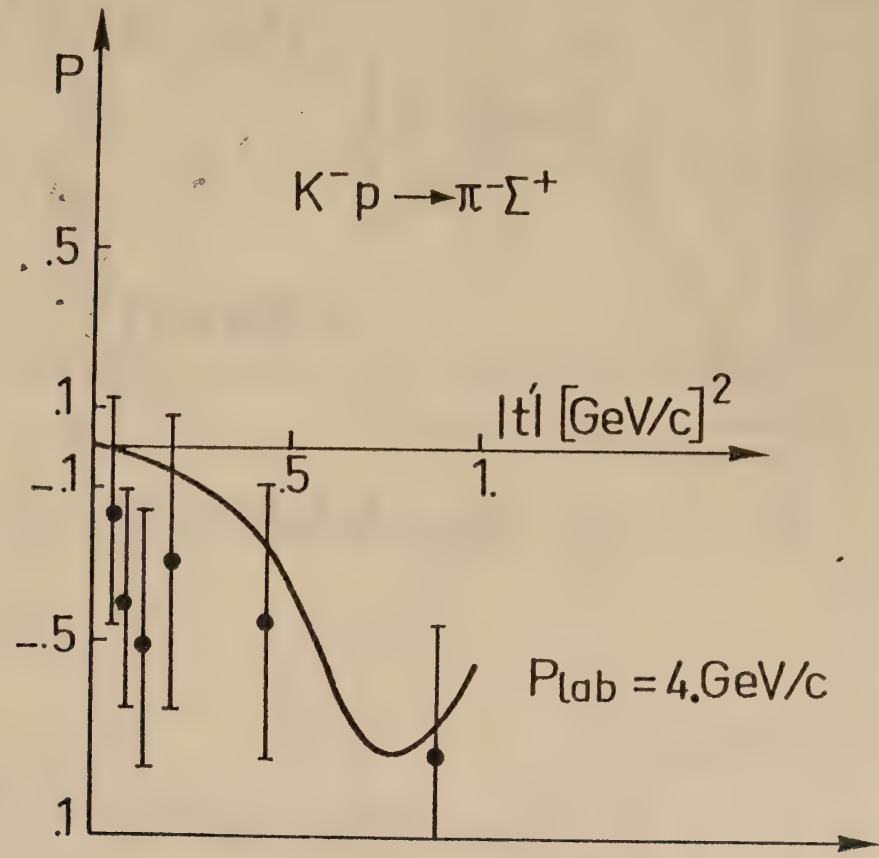

P
$K^- p \rightarrow \pi^- \Sigma^+$
.5
.1
-.1
-.5
.1
$|t'|$ [GeV/c]^2
.5
1.
P_{lab} = 4.GeV/c

fig 4-a

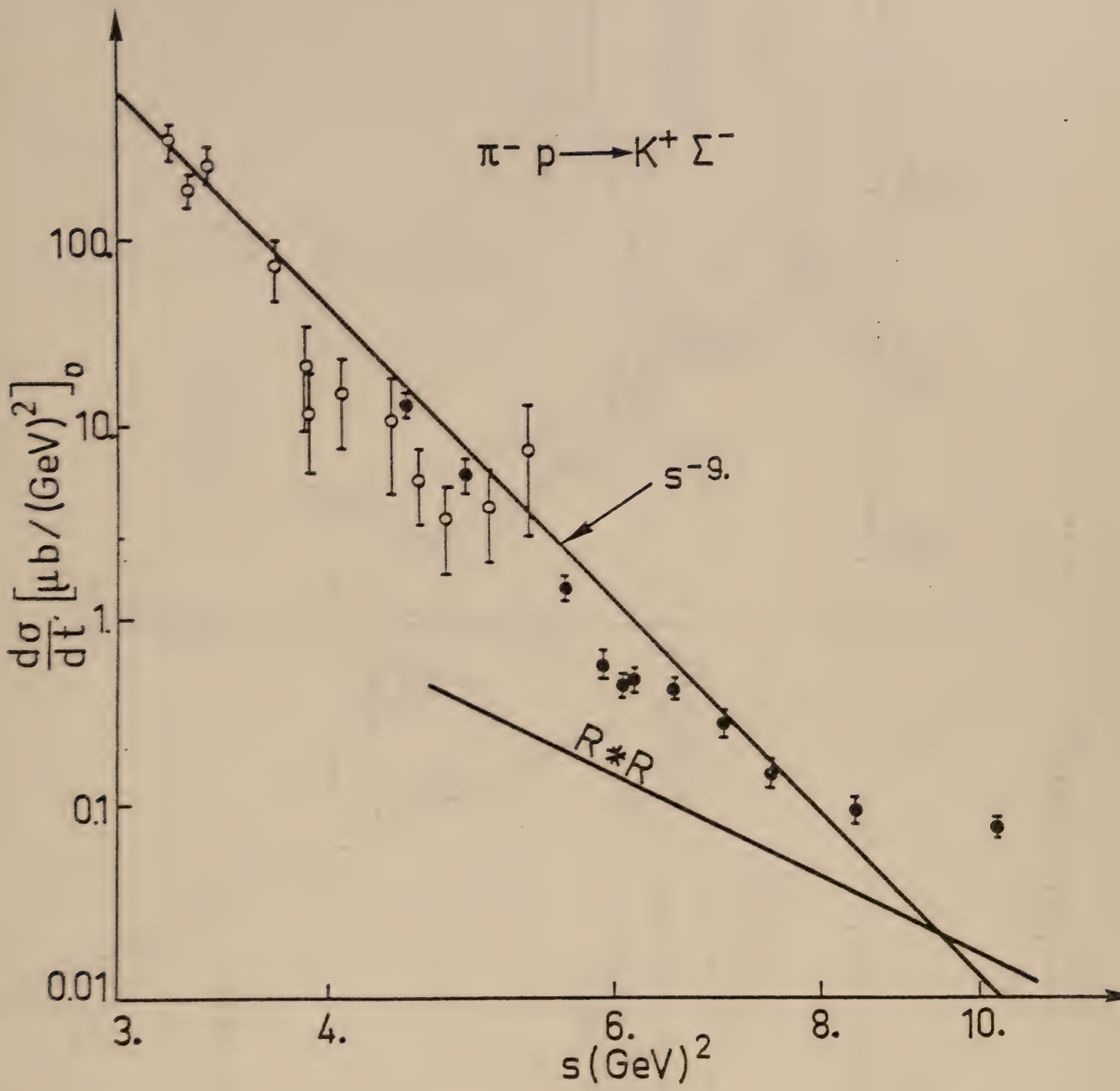
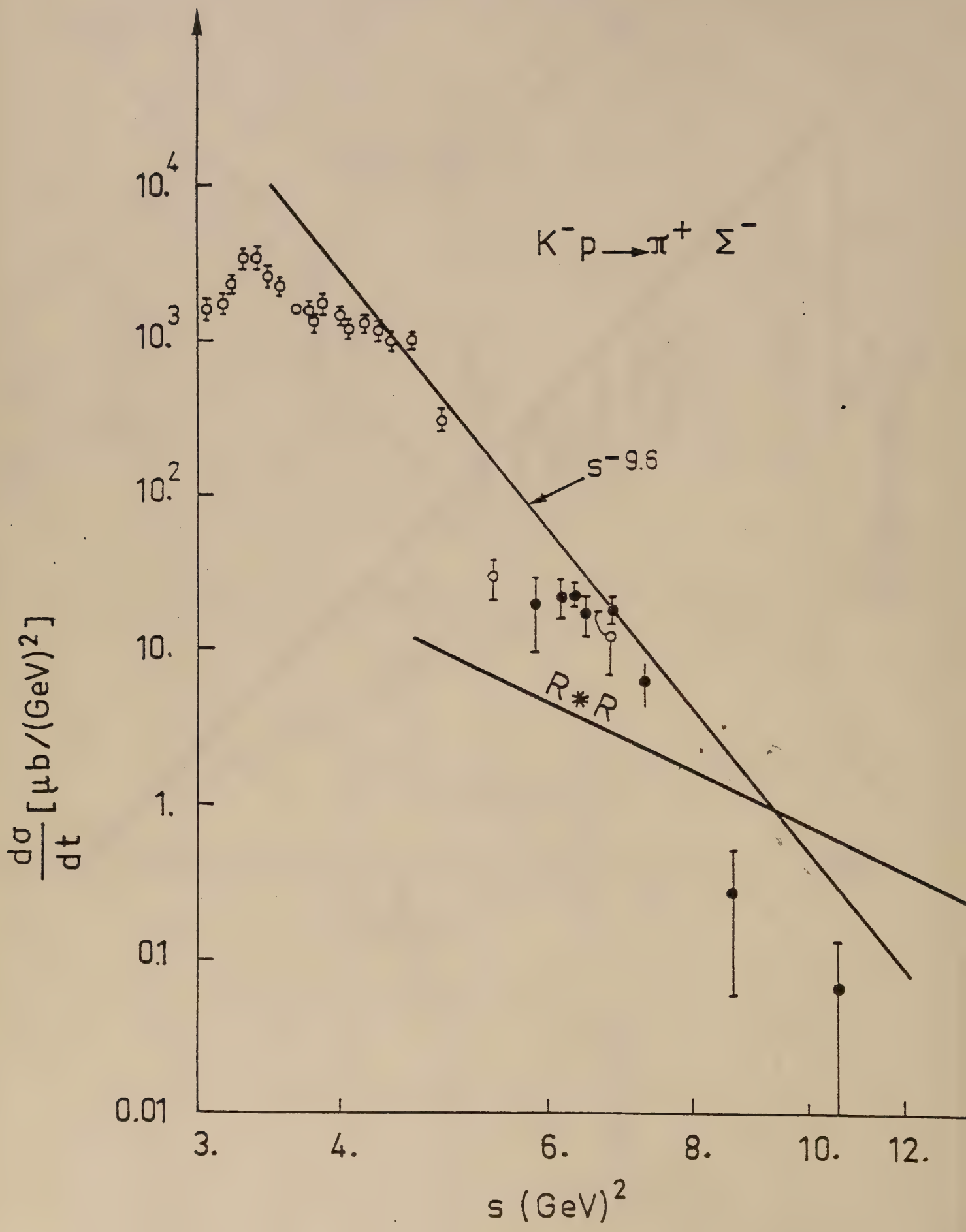

$\pi^- p \longrightarrow K^+ \Sigma^-$
100.
10.
1.
0.1
0.01
$\frac{d\sigma}{dt}\left[\mu b/(GeV)^2\right]_0$
$s^{-9.}$
R*R
3.
4.
6.
8.
10.
$s\,(GeV)^2$

fig 4-b


$K^- p \rightarrow \pi^+ \Sigma^-$
$s^{-9.6}$
$R*R$
$\frac{d\sigma}{dt}$ [μb/(GeV)²]
10.⁴
10.³
10.²
10.
1.
0.1
0.01
3.
4.
6.
8.
10.
12.
s (GeV)²

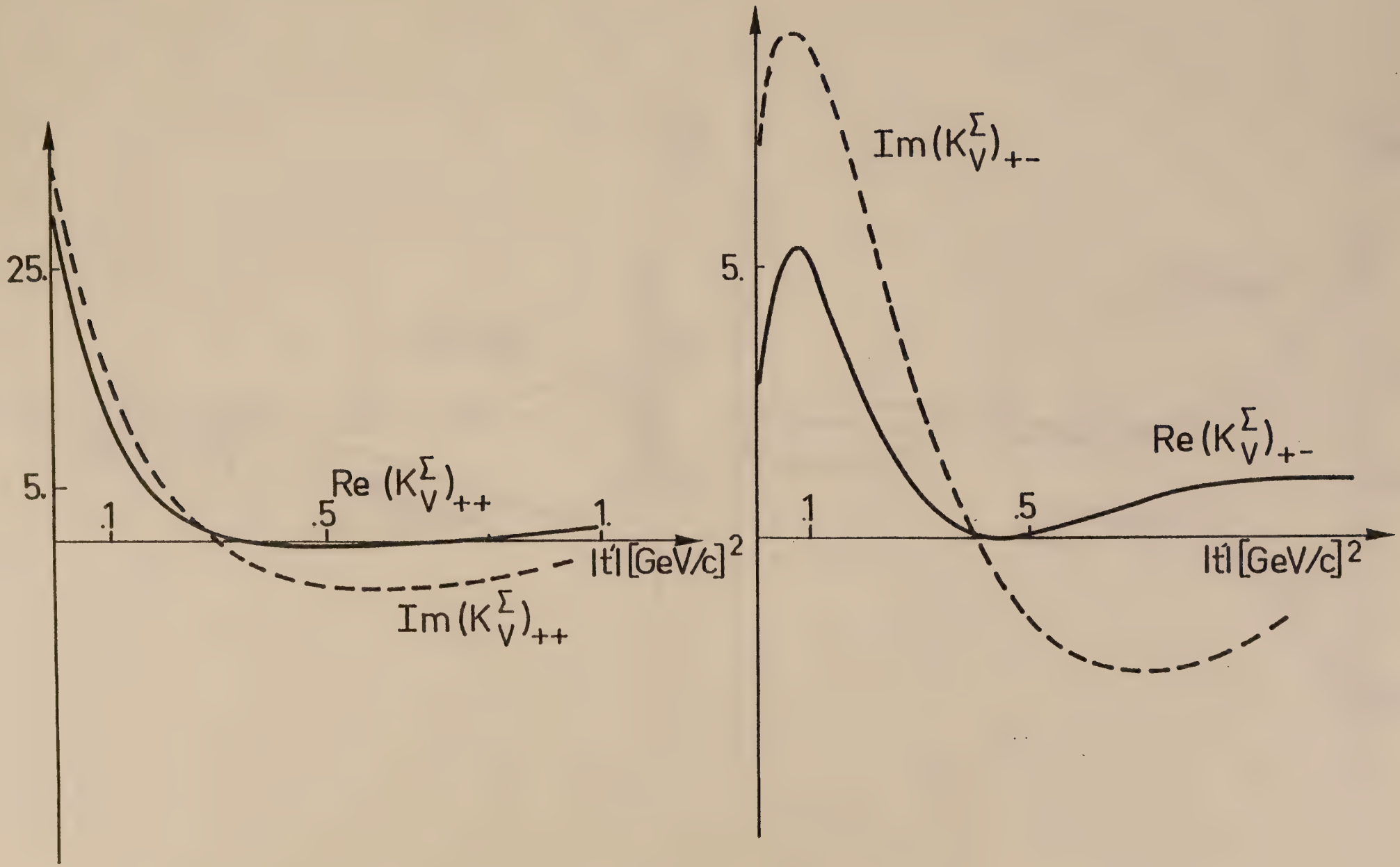

25.
5.
.1
.5
1.
Re $(K_V^{\Sigma})_{++}$
Im $(K_V^{\Sigma})_{++}$
$|t'|$ [GeV/c]^2
5.
1
.5
Im $(K_V^{\Sigma})_{+-}$
Re $(K_V^{\Sigma})_{+-}$
$|t'|$ [GeV/c]^2

Fig 5-a

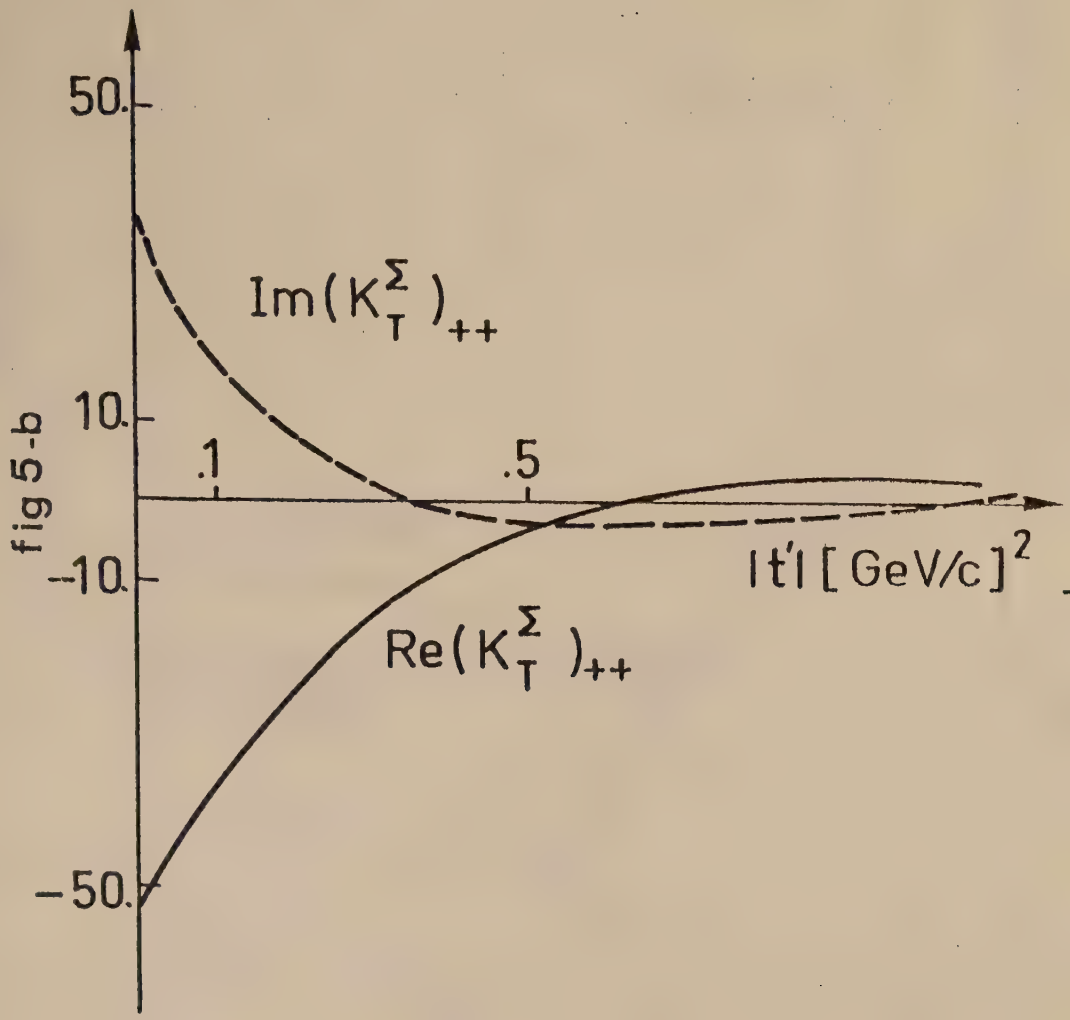
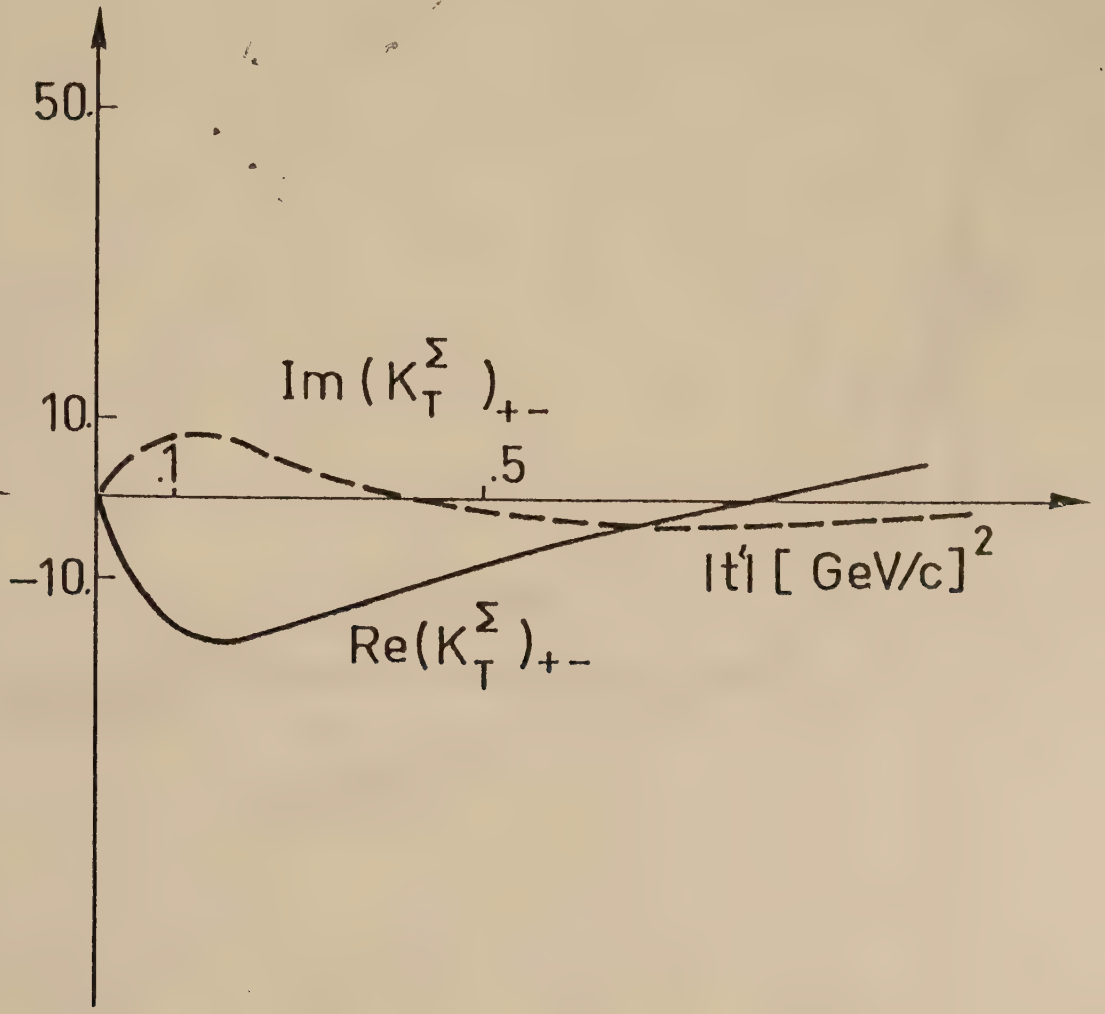

50.
10.
-10.
-50.
.1
.5
$Im(K_T^{\Sigma})_{++}$
$Re(K_T^{\Sigma})_{++}$
$|t'|$ [GeV/c]^2

fig 5-b


50.
10.
-10.
.1
.5
$Im(K_T^{\Sigma})_{+-}$
$Re(K_T^{\Sigma})_{+-}$
$|t'|$ [GeV/c]^2

2ème PARTIE

ANALYSE EN AMPLITUDES DES REACTIONS

MESON-BARYON A 6 GEV/c

:-:-:-:

INTRODUCTION

La deuxième partie de ce travail est consacrée à des analyses en amplitudes à énergie fixée. (Dans tous les cas les données expérimentales utilisées correspondent à une impulsion de 6 GeV/c dans le laboratoire). L'attitude adoptée est très différente de celle de la première partie, dans laquelle nous avons construit un modèle "ab initio" et en avons exploré toutes les prédictions expérimentales. Ici au contraire le problème posé est la résolution des équations reliant les amplitudes de diffusion aux observables mesurées expérimentalement. Le point de départ est le résultat expérimental, nous n'introduisons des hypothèses théoriques qu'afin de rendre possible la résolution des équations. Ainsi dans la diffusion $\pi N \to \pi N$ où un ensemble complet d'observables a été mesuré, il est possible de faire une analyse en amplitudes de façon intrinsèque, sans aucune hypothèse théorique. Ces résultats ont donné des informations très intéressantes sur les amplitudes d'échange. En particulier, elles confirment les constatations des modèles à savoir :

- i) L'amplitude d'isospin 0 dans la voie t, $n=0$ [$n=|\lambda_3-\lambda_4-\lambda_1+\lambda_2|$ est le bilan d'hélicité dans la voie s] est très fortement dominante et principalement imaginaire[†]. La composante $n = 1$ de cette amplitude est très faible.

- ii) L'amplitude d'échange d'isospin 1, $n = 0$ a un zéro périphérique dans sa partie imaginaire et la composante $n = 1$, qui est plus couplée à cette voie a une structure compatible avec celle d'un pôle de Regge.

La situation est moins favorable dans les autres réactions où seules quelques observables ont été mesurées. Il est alors nécessaire de compléter le système en introduisant des hypothèses auxiliaires. Dans le chapitre suivant, nous présentons une étude des réactions $\pi^- p \to \eta n$ et $K^- p \to \overline{K}^0 n$ qui permet d'obtenir les amplitudes d'échange des nombres quantiques du A_2

[†] Les phases peuvent être obtenues par l'étude de l'interférence Coulombienne.

(nous dirons échange du A_2 par abus de langage). Cette étude est intéressante comme test de la dégénérescence d'échange, qui devrait exister entre le ρ et le A_2 .

Pour réaliser cette analyse, nous introduisons deux hypothèses théoriques, hypothèses que nous utiliserons dans toutes les analyses ultérieures.

(a) La symétrie SU(3) est exacte. Nous pouvons ainsi relier des réactions très différentes et disposer de plus d'information expérimentale.

(b) L'amplitude d'échange du ρ (n=1) a une phase de Regge $i \exp(-\frac{i\pi\alpha(t)}{2})$. Cette hypothèse est nécessaire afin de fixer les phases absolues des amplitudes, l'analyse de la réaction πN étant faite à une phase globale près. Nous obtenons alors les informations suivantes sur l'échange du A_2 jusqu'à des valeurs de $|t|$ de $.7(GeV/c)^2$.

- iii) L'amplitude $A_2(n=1)$ a un comportement de pôle de Regge de signature positive, dégénéré d'échange avec le $\rho(n=1)$. Ce qui confirme l'idée que l'absorption joue peu sur les composantes (n=1) et modifie très peu la nature du pôle de Regge.

- iv) L'amplitude A_2 (n=0) a une structure assez particulière. Il semble que la partie imaginaire ait un zéro double et la partie réelle un zéro simple à des valeurs de t voisines de $|t| \simeq .6\ (GeV/c)^2$. Des résultats plus précis présentés au chapitre VII confirment ces observations. Un tel comportement est bien décrit par la forme paramétrique suivante :

$$A_2(n=0) \sim A\ \alpha(t)\ e^{-i\frac{\pi\alpha}{2}(t)} \quad .$$

Il convient de souligner ici que des études à l'aide de règles de somme à énergie finie donnent une structure comparable pour le $f_o(n=0)$ - de même que le modèle RPR au chapitre II, figure 3-19-. Les amplitudes du ρ et du A_2 étant connues, nous étudions les réactions $K^{\pm}p \to K^{\pm}p$ et la régénération du K^o : $K_L^o p \to K_s^o p$ (chapitre V). Une étude très détaillée des mesures expérimentales nous conduit à poser :

$$\rho(n = 0,1) = A_n\ \omega(n = 0,1) \ ,$$

où la constante de proportionnalité A_n est indépendante du transfert. Ce résultat est très important car il suggère que les échanges vecteurs ($J^P = 1^-$) sont étroitement reliés. Cette relation nous permet alors d'extraire l'amplitude d'échange du Pomeron et du f_o, c'est-à-dire d'isospin O et de signature positive M_n^{O+}.

- v) M_o^{O+} dans les réactions KN est beaucoup plus importante que M_1^{O+} et pratiquement imaginaire pure. Toutefois, le rapport M_1^{O+}/M_o^{O+} est plus grand que dans πN.

Dans le chapitre VI, où nous étudions les réactions $\pi^+ p \to \pi^o(\eta)\ \Delta^{++}$, $K^+ p \to K^o \Delta^{++}$ et $K^- n \to \bar{K}^o \Delta^-$ nous devons faire plus d'hypothèses. En effet SU(3) ne donne aucune relation entre les réactions ci-dessus et celles précédemment étudiées. Nous supposons que le ρ et le A_2 qui contribuent ici sont identiques à ceux qui apparaissent dans $\pi N \to \pi(\eta)N$ à des couplages près (indépendants de t). Cette hypothèse (qui est bien vérifiée) nous permet d'obtenir les amplitudes $\rho(n=2)$ et $A_2(n=2)$. Les incertitudes sont néanmoins trop importantes pour permettre de tirer des conclusions définitives. L'exactitude de notre hypothèse signifie que la structure en t d'une amplitude d'échange de nombres quantiques bien définies dépend de n et non des vertex dans la voie t. Cette constatation est très proche de l'image géométrique de la diffusion où pour un n donné, nous avons $\mathrm{Im}\, M_n(s,t) \sim J_n(R\sqrt{-t})$. Nous avons déjà souligné que les amplitudes d'échange du ρ et du ω étaient semblables, de même que le f_o est à rapprocher du A_2. Les amplitudes paraissent pouvoir être regroupées en deux classes : les signatures négatives (vecteurs 1^-) et les signatures positives (tenseurs 2^+).

Dans le dernier chapitre, nous présentons précisément un test de cette universalité vecteur-tenseur étendue au $K^*(890)$ et au $K^{**}(1420)$. En utilisant simplement le ρ et le A_2 déjà déterminés on reproduit de façon très correcte les données expérimentales dans les réactions d'échange d'hypercharge. Le ρ et le A_2 tiennent lieu de "prototypes" des échanges vecteur et tenseur.

On voit comment l'approche des analyses en amplitude, en mettant au jour des régularités dans les amplitudes, permet peu à peu de faire émerger une systématique simple. La systématique que nous proposons est est résumée dans la table I du dernier chapitre.

On peut se poser la question de la dépendance en énergie. Ainsi H. Navelet et P. Stevens ont utilisé des formes paramétriques ayant les propriétés énoncées, ils ont pu obtenir une très bonne description des réactions d'échange d'hypercharge entre 4 GeV/c et 17 GeV/c. Cette systématique permet de simplifier considérablement la description des réactions méson-baryon, car elle implique quantité de relations entre réactions. Le problème reste cependant ouvert quant à l'explication théorique d'une telle classification. En langage du paramètre d'impact il faut comprendre pourquoi lorsque le potentiel de la voie t est du type vecteur (resp. tenseur) le hadron apparaît comme un anneau (resp. un disque) de matière hadronique.

A ce point il faut signaler une tentative théorique de description globale des amplitudes par la b-universalité, à savoir que les fonctions de profil sont de la forme :

$$f_n(b,s) \sim b^n f_o(b,s) \quad .$$

Cette approche semble donner de bons résultats pour le ρ et le A_2 . Le problème paraît en fait plus lié à la dynamique interne et à la structure des hadrons eux-mêmes qu'au modèle d'échange.

CHAPITRE IV

AMPLITUDE ANALYSIS OF THE A_2 EXCHANGE AT 6 GEV/C

by

G. GIRARDI, C. GODRECHE and H. NAVELET

AMPLITUDE ANALYSIS OF THE A_2 EXCHANGE AT 6 GEV/C

by

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ABSTRACT

An analysis of experimental data on $\pi^- p \to \eta n$ and $K^- p \to \bar{K}^0 n$ reactions is performed which enablesus to exhibit the A_2 amplitudes. There are strong evidences that the A_2 spin flip amplitude behaves like a Regge amplitude exchange degenerate with the ρ spin flip one.

In the recent years a lot of experimental data became available which enable to test high energy models by performing amplitude analysis.

For instance in πN scattering at 6 GeV/c the measurement of the spin rotation parameters (A and R) of elastic reactions[1] provided a complete set of observables and has allowed a model independent amplitude analysis. Several works were published which determined the four amplitudes of these reactions, up to an overall phase[2-5]. In KN scattering a complete set of experimental data is not yet available ; nevertheless it is possible by using quite a few reasonable assumptions to extract the A_2 amplitudes(**), thanks to the recent measurement of the polarisation of the recoil neutron in $K^- p \to \bar{K}^0 n$ at 8 GeV/c[6].

(*) L P T H E - Université Paris VII

(**) A_2 refers to the exchange in the t-channel of the quantum numbers of A_2 and not to A_2 as a Regge pole. This is only a labelling of t-channel quantum numbers.

Our conventions are the followings :

$$\frac{d\sigma}{dt} = |M_+|^2 + |M_-|^2$$
$$P\frac{d\sigma}{dt} = -2\,\mathrm{Im}\,(M_+ M_-^*) \quad , \tag{1}$$

where $M_\pm$ are the s-channel helicity amplitudes and the subscripts $\pm$ refer to the non flip and the flip amplitudes.

All along this work we shall assume that SU(3) is exact, except for the Pomeron exchange. With this hypothesis one can write :

$$M_i(\pi^\pm p \to \pi^\pm p) = (P + f_o)_i = \rho_i \tag{2a}$$
$$M_i(\pi^- p \to \pi^o n) = -\sqrt{2}\,\rho_i \tag{2b}$$
$$M_i(\pi^- p \to \eta n) = \sqrt{\frac{2}{3}}\,(A_2)_i \tag{2c}$$
$$M_i(K^- p \to \bar{K}^o n) = (\rho - A_2)_i \tag{2d}$$
$$M_i(K^+ n \to K^o p) = -(\rho + A_2)_i \quad , \tag{2e}$$

where $i = (\pm)$, P represents the vacuum quantum numbers, the other letters are standard notations.

Exact SU(3) implies a relation between cross sections of reactions (2b-d) which is well satisfied by the data :

$$\frac{d\sigma}{dt}(K^- p \to \bar{K}^o n) + \frac{d\sigma}{dt}(K^+ n \to K^o p) = 3\,\frac{d\sigma}{dt}(\pi^- p \to \eta n) + \frac{d\sigma}{dt}(\pi^- p \to \pi^o n) \quad . \tag{3a}$$

This relation can be extended to the polarisations, namely.

$$P\,\frac{d\sigma}{dt}(K^- p \to \bar{K}^o n) + P\,\frac{d\sigma}{dt}(K^+ n \to K^o p) = 3P\,\frac{d\sigma}{dt}(\pi^- p \to \eta n) + P\,\frac{d\sigma}{dt}(\pi^- p \to \pi^o n). \tag{3b}$$

Since we have assumed that SU(3) holds, both $\frac{d\sigma}{dt}(K^+ n \to K^o p)$ and $P(K^+ n \to K^o p)$ will be determined by (3a-b) once we know the data of the 3 other reactions. We shall use the following experimental data :

$$\frac{3}{2}\frac{d\sigma}{dt}(\pi^- p \to \eta n) = d\sigma^e = |A_{2+}|^2 + |A_{2-}|^2 \quad \text{at} \quad 5.9 \text{ GeV/c}^{[7]} \qquad (4a)$$

$$P^e\, d\sigma^e = -2\,\mathrm{Im}\,[A_{2+} \cdot A^*_{2-}] \quad \text{at} \quad 4.9 \text{ and } 7.85 \text{ GeV/c}^{[8]} \qquad (4b)$$

$$\frac{d\sigma}{dt}(K^- p \to \bar{K}^o n) = d\sigma^k = |(\rho\text{-}A_2)_+|^2 + |(\rho\text{-}A_2)_-|^2 \text{ at 5 and 7 GeV/c}^{[9]} \qquad (4c)$$

$$P^k\, d\sigma^k = -2\,\mathrm{Im}\,[(\rho\text{-}A_2)_+ (\rho\text{-}A_2)^*_-] \quad \text{at} \quad 8 \text{ GeV/c}^{[6]}\,. \qquad (4d)$$

To obtain the data set at 6 GeV/c, we have interpolated for P^e and $d\sigma^k$ between the 2 sets of measurements, for P^k we have assumed a smooth energy dependence, the corresponding values are given in table I. As aforementioned, through amplitude analysis, the amplitudes of πN reactions are known up to an overall phase. Nevertheless in all these analyses[2-5] the ρ exchange spin flip amplitude, ρ_-, seems to behave like a pure Regge pole exchange amplitude. Hence it is reasonable to assume that ρ_- has effectively the Regge phase, this assumption removes the overall phase ambiguity. In our analysis we have chosen the results of ref.[2] in which the modulus of ρ_- exhibits a minimum around $t = -.75\ (\text{GeV/c})^2$. Then we parametrise ρ_- with the following form :

$$\rho_-(t) = |\rho_-(t)|\, e^{-\frac{i\pi}{2}(1-\alpha(t))}\,, \qquad (5)$$

where $\alpha(t) = .56\,(1 + \frac{4t}{3})$ labels the ρ trajectory which has the correct intercept and reproduces the observed dip. The s dependence is irrelevant since we work at fixed energy.

Now from equations (4) we are able to extract 4 numbers $\mathrm{Im}(A_2)_\pm$ and $\mathrm{Re}(A_2)_\pm$. However the following trick makes the resolution easier : we use the transversity amplitudes namely :

$$(A'_2)_\pm = \frac{1}{\sqrt{2}}\,[(A_2)_+ \pm i(A_2)_-]$$

$$= |A'_2|_\pm \exp i(\varphi'_{A_2})_\pm \qquad (6a)$$

and

$$(\rho')_\pm = \frac{1}{\sqrt{2}}\,[(\rho)_+ \pm i(\rho)_-] = |(\rho')_\pm| \exp i\,(\varphi'_\rho)_\pm\,, \qquad (6b)$$

here the subscript $(\pm)$ only refers to the chosen sign in the sum of the right hand member of (6).

Using (6-ab) the set of equations (4) becomes

$$d\sigma^e = |(A_2')_+|^2 + |(A_2')_-|^2 \qquad (7a)$$

$$p^e d\sigma^e = |(A_2')_-|^2 - |(A_2')_+|^2 \qquad (7b)$$

$$d\sigma^k = |(\rho'-A_2')_+|^2 + |(\rho'-A_2')_-|^2 \qquad (7c)$$

$$p^k d\sigma^k = |(\rho'-A_2')_-|^2 - |(\rho'-A_2')_+|^2 \quad , \qquad (7d)$$

the solutions of this system are

$$|(A_2')_\pm|^2 = \tfrac{1}{2}\, d\sigma^e\, (1 \mp p^e) \qquad (8)$$

$$(\varphi'_{A_2})_\pm = (\varphi'_\rho)_\pm + \varepsilon \operatorname{Arc\,cos} C_\pm \qquad (9)$$

where $C_\pm = \dfrac{1}{2\sqrt{2}} \dfrac{[d\sigma^e(1 \mp p^e) + 2|\rho'_\pm|^2 - d\sigma^k(1 \mp p^k)]}{|\rho'_\pm| \sqrt{d\sigma^e\,(1 \mp p^e)}}$,

and the $\varepsilon = \pm 1$ ambiguity is due to the Arcos $C_\pm$ function.

Thus we have 2 solutions for $(A_2')_+$ and 2 solutions for $(A_2')_-$; going back to the helicity amplitudes $(A_2)_+$ and $(A_2)_-$ we obtain 4 solutions for each of them. One can try to follow each solution by continuity, but, whenever $C_\pm$ goes through the value $+1$, one looses the track and one can jump from one solution to another. In order to select a solution, one has to use some physical considerations.

(a) The spin flip amplitudes $(A_2)_-$ must vanish at $t = 0$, but not $(A_2)_+$, this eliminates near the forward direction 2 solutions (noted 2 and 3) where the roles of $(M)_+$ and $(M)_-$ are exchanged. This confusion is only due to the fact that in (1) we have not imposed the kinematical singularity of $(M)_-$.

(b) Between 5 and 15 GeV/c , $\sigma_{tot}(K^+p)$ is flat, thus we expect that $Im\,(A_2)_+ \sim -\,Im(\rho)_+$ at very low t , the solution n°4 does not fulfill this condition. Finally we are left with only one acceptable solution (n°1) close to $t = 0$. Furthermore this solution enjoys two others properties.

i) $|A_2|_- > |A|_+$ near $t = 0$ which is consistent with the shape of $d\sigma^e$ and $d\sigma^k$ (turn-over point) which indicates that $(M)_-$ dominates over $(M)_+$.

ii) $Im(\rho)_- \simeq -\,Im(A_2)_-$ up to $t = -.5\ (GeV/c)^2$. This is in agreement with a very important theoretical consequence of duality namely : exchange degeneracy. Indeed if $(A_2)_-$ and $(\rho)_-$, which has the Regge phase, are exchange degenerate, this yields a small difference between $\frac{d\sigma}{dt}(K^+n \to K^op)$ and $\frac{d\sigma}{dt}(K^-p \to \bar{K}^on)$ [flip dominance (i)] almost constant in t , property which is observed. In fact solution 1 is not quite satisfactory over all the t range we have investigated $[0 < |t| < .7\ (GeV/c)^2]$: after a minimum around $t = -.45\ (GeV/c)^2$, $|A_2|_+(1)$ increases drastically and even becomes larger than $|A_2|_-(1)$. However close to $t = -.5\ (GeV/c)^2$, C_- goes through $+1$ and then there is a bifurcation point where solutions 1 and 2 coincide [fig.1,2] . Thanks to this ambiguity we can preserve the property (i) all along the t range by passing continuously from solution 1 to solution 2 . Nevertheless there are no rigourous arguments to choose one path rather than the other at this bifurcation point.

The final results are given in table II and in figure 3 with their error bars. The uncertainties were calculated by assuming that all the errors were uncorellated, then they are overestimated.

One can summarize the main features of the results as follows :

- With only a few hypotheses : SU(3) symmetry and the Regge phase for the ρ spin flip we were able to gain important information on the A_2 exchange.

The spin flip amplitude $(A_2)_-$ is well determined and behaves like a pure Regge pole of positive signature. Moreover there is strong evidence that it is exchange degenerate with the ρ spin flip up to $t = -.5$ GeV/c , see fig.(3c-d) over this t range the corresponding phase (φ_{A_2}) (fig.2b) may be well described by a linear Regge trajectory

$$\alpha_{A_2}(t) = .55 + .85\,t \quad . \qquad (10)$$

- Beyond $-.5\ (GeV/c)^2$ in transfert momentum we found a change in the trajectory indicating that in this region the Regge behaviour is strongly violated.

As far as the non flip amplitude $(A_2)_+$ is concerned it is badly determined since it is small as compared to (A_2). Anyhow one can make some remarks.

- $(A_2)_+$ does not behave like a Regge pole.

- At very low t $Im\ (A_2)_+ \sim -\ Im(\rho)_+$, but whereas $Im\ \rho_+$ vanishes at $t \sim -.2\ (GeV/c)^2$. $Im\ (A_2)_+$ seems - within the errors bars - to have a double zero around $t \sim -.6\ (GeV/c)^2$ (fig.2-ab).

Finally through relation (3-ab) we give the predictions obtained for $\frac{d\sigma}{dt}(K^+n \to K^op)$ and $P(K^+n \to K^op)$ by using the values of the observables of the 3 other reactions we have taken as input (fig.4-ab).

This analysis seems to give a good evidence for exchange degeneracy between the ρ and A_2 spin flip amplitudes. As far as high energy models are concerned this confirms that absorption - if there is - has to be much weaker in spin flip amplitudes than in non flip amplitudes which are quite different from simple Regge poles. Using data on KN elastic scattering and regeneration we are investigating the ω and - if possible - the f_o exchange amplitudes.

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TABLE I

$-t(GeV)^2$	$d\sigma^e$ $mb/(GeV)^2$	p^e %	$d\sigma^k$ $mb/(GeV)^2$	p^k %
.04	.18	- .12	.32	- .02
.12	.21	- .09	.34	- .05
. 2	.20	- .04	.26	- .07
. 3	.14	.04	.16	- .09
. 4	.09	.11	.096	- .2
. 5	.06	.12	.056	- .27
. 6	.04	.1	.035	- .33
. 7	.027	.09	.023	- .25

TABLE II

$-t$ $(GeV)^2$	Re A_+ $(mb)^{1/2}(GeV)^{-1}$	Im A_+ $(mb)^{1/2}(GeV)^{-1}$	Re A_- $(mb)^{1/2}(GeV)^{-1}$	Im A_- $(mb)^{1/2}(GeV)^{-1}$
.04	.190 ± .125	-.132 ± .090	.263 ± .101	-.240 ± .123
.12	.087 ± .068	-.044 ± .078	.346 ± .064	-.283 ± .040
. 2	.067 ± .050	-.022 ± .065	.373 ± .045	-.234 ± .033
. 3	.053 ± .063	-.036 ± .067	.324 ± .058	-.168 ± .053
. 4	.033 ± .073	-.031 ± .036	.279 ± .036	-.107 ± .056
. 5	-.020 ± .085	-.017 ± .026	.242 ± .038	.030 ± .132
. 6	.009 ± .045	-.005 ± .061	.085 ± .070	.180 ± .092
. 7	.015 ± .055	-.038 ± .043	-.031 ± .046	.156 ± .079

FIGURES CAPTIONS

Fig.1 : Moduli of $(A_2)_+$ and $(A_2)_-$, solution n° 2 is drawn only after the bifurcation point.

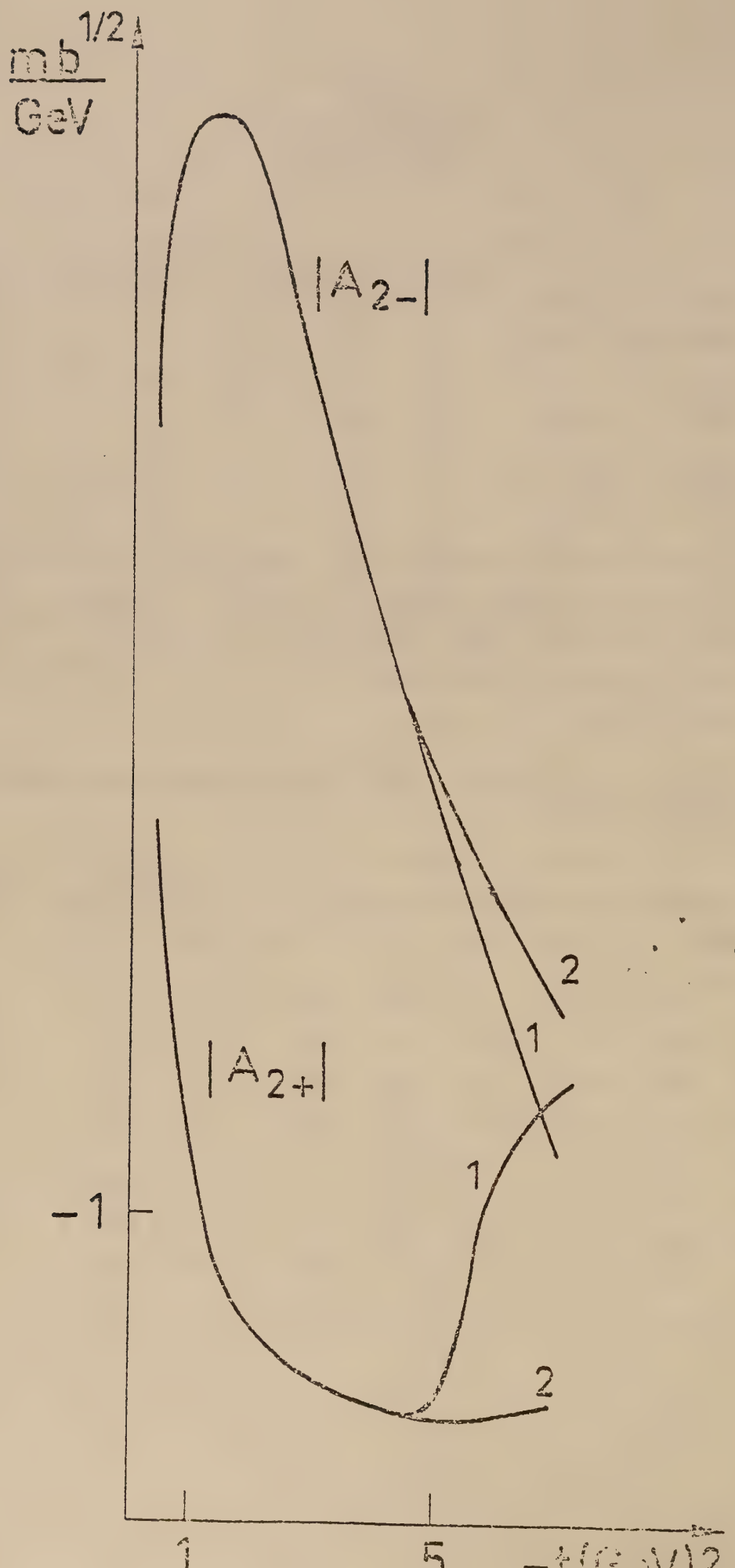
Fig.2a : Phase of $(A_2)_+$

Fig.2b : Phase of $(A_2)_-$

Fig.3 : Real part and Im part of $(A_2)_+$ and $(A_2)_-$. The heavy lines give the values obtained by using data of table I. Upper and lower solid lines are the error bars. The dashed lines are the ρ amplitudes of ref.[2] given for comparison, where the sign is changed for the imaginary parts.

Fig.4a : Prediction of $\frac{d\sigma}{dt}(K^+n \rightarrow K^\circ p)$ using relation (3a) data are from ref.[10]

Fig.4b : Prediction of the polarization in $K^+n \rightarrow K^\circ p$ reaction using relation (3b)


mb 1/2
GeV
|A2−|
|A2+|
2
1
1
2
−1
.1
.5
−t(GeV)2

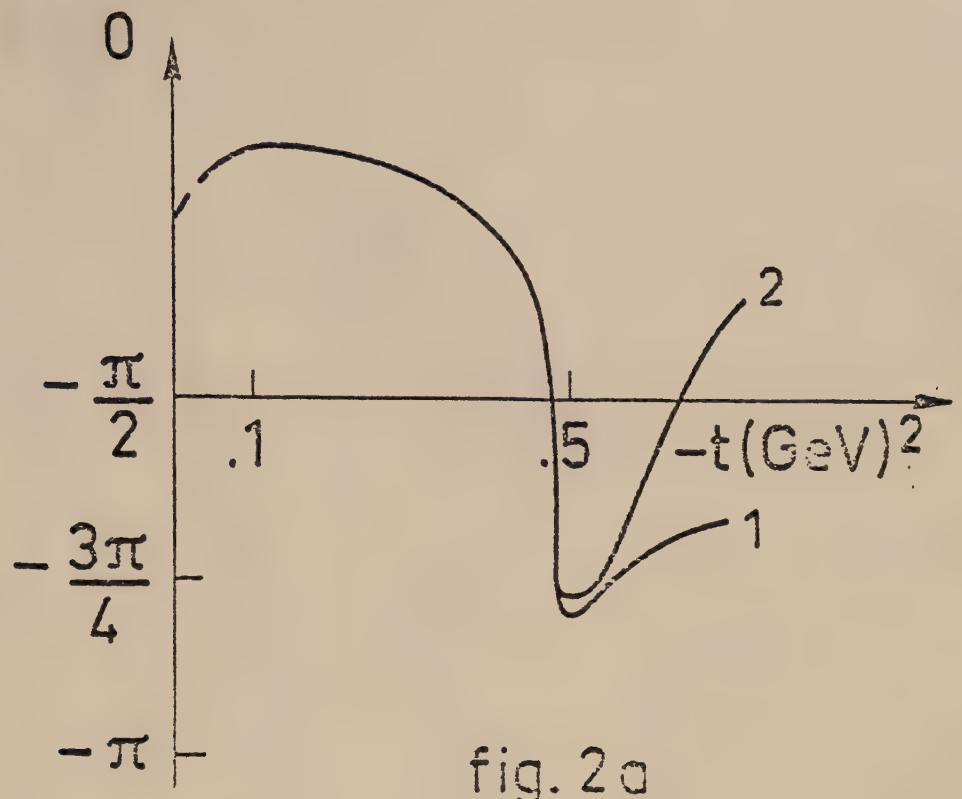


fig. 2a
$(\varphi A_2)+$

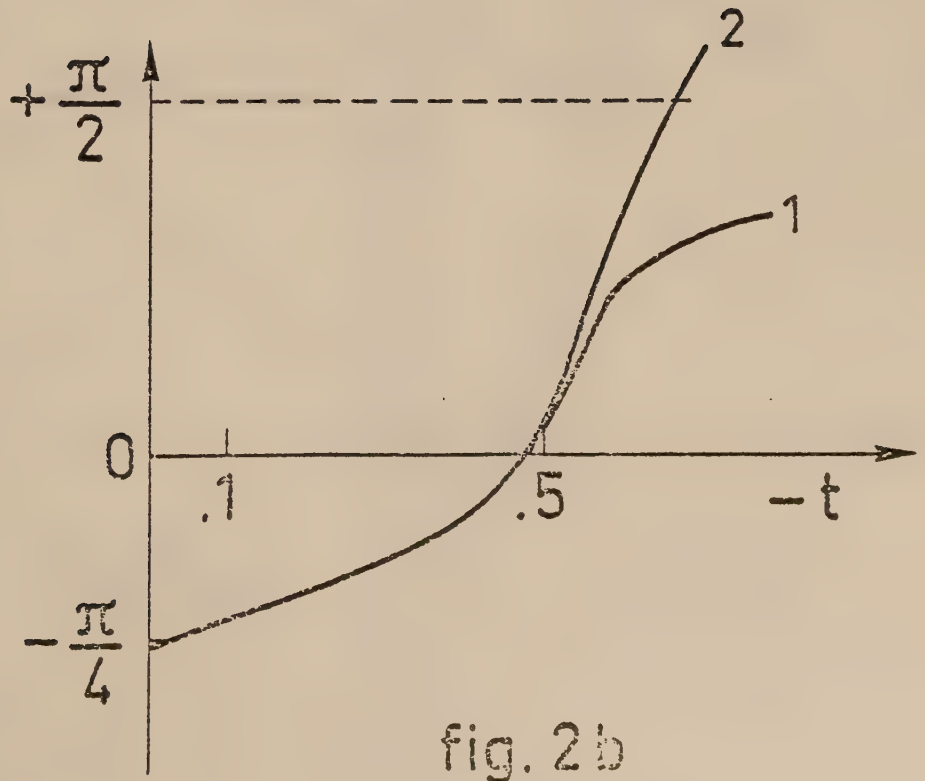
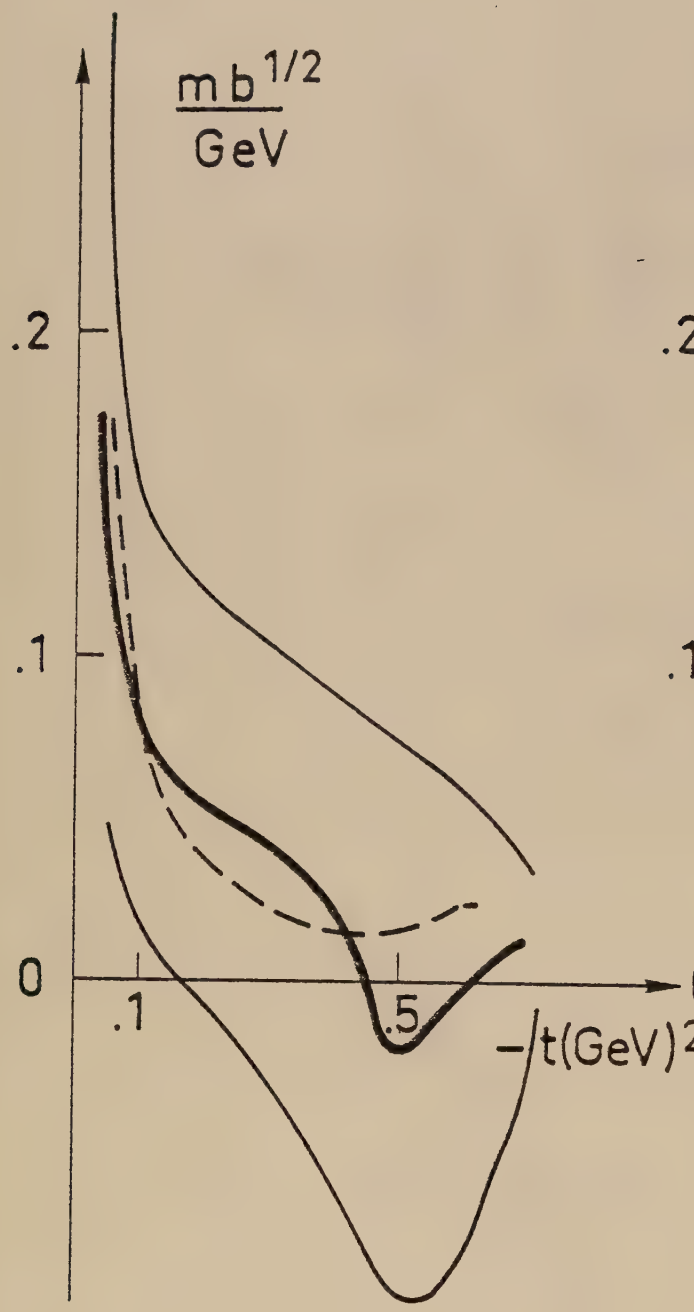
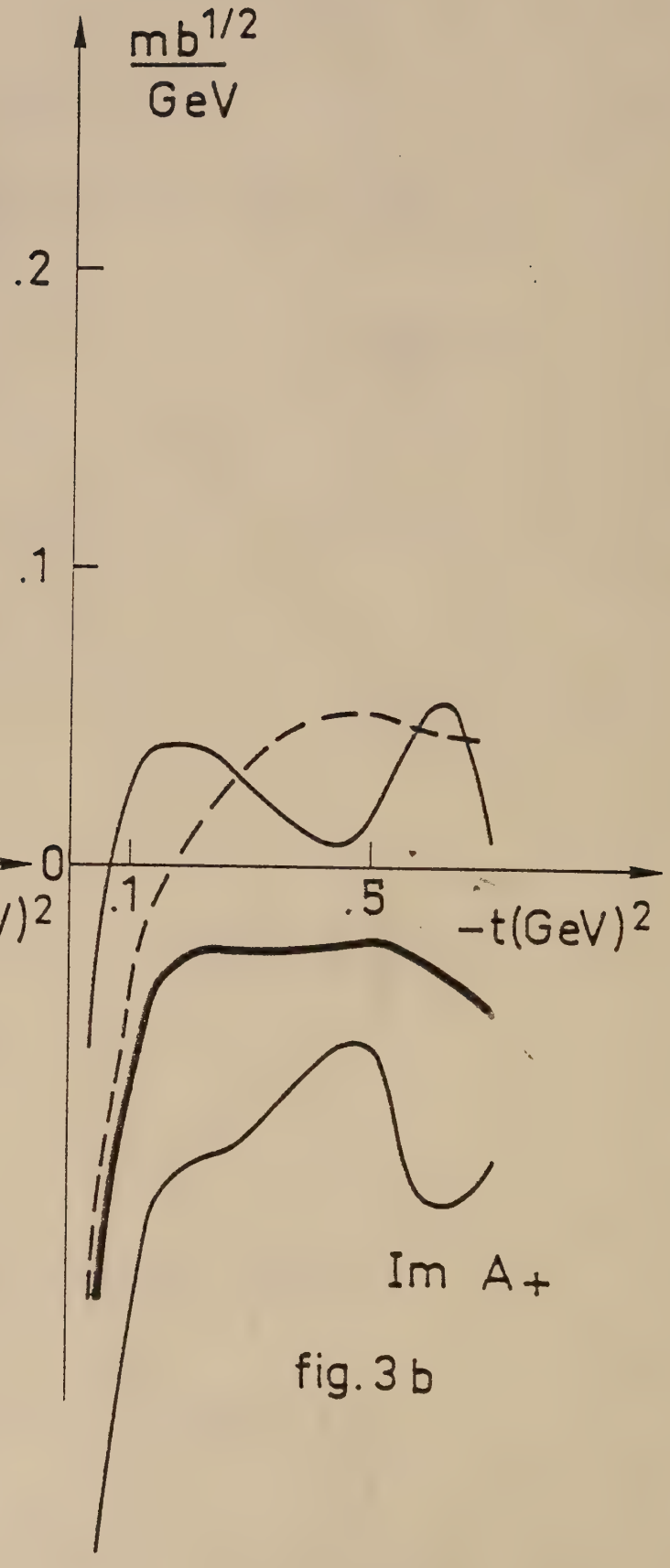


fig. 2b
$(\varphi A_2)-$



Re (A_+)

fig. 3 a



Im A_+

fig. 3 b

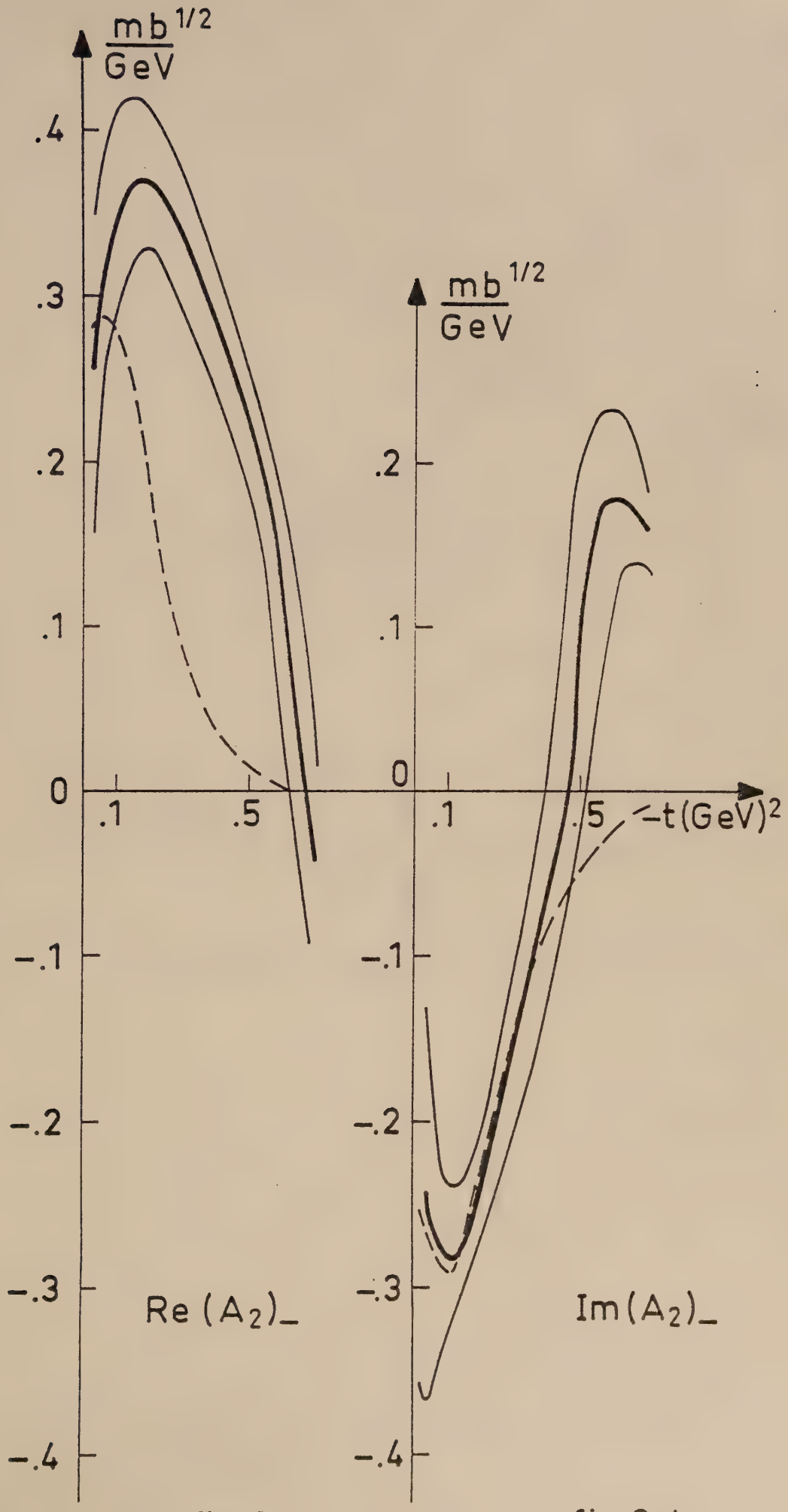


fig. 3 c

fig. 3 d

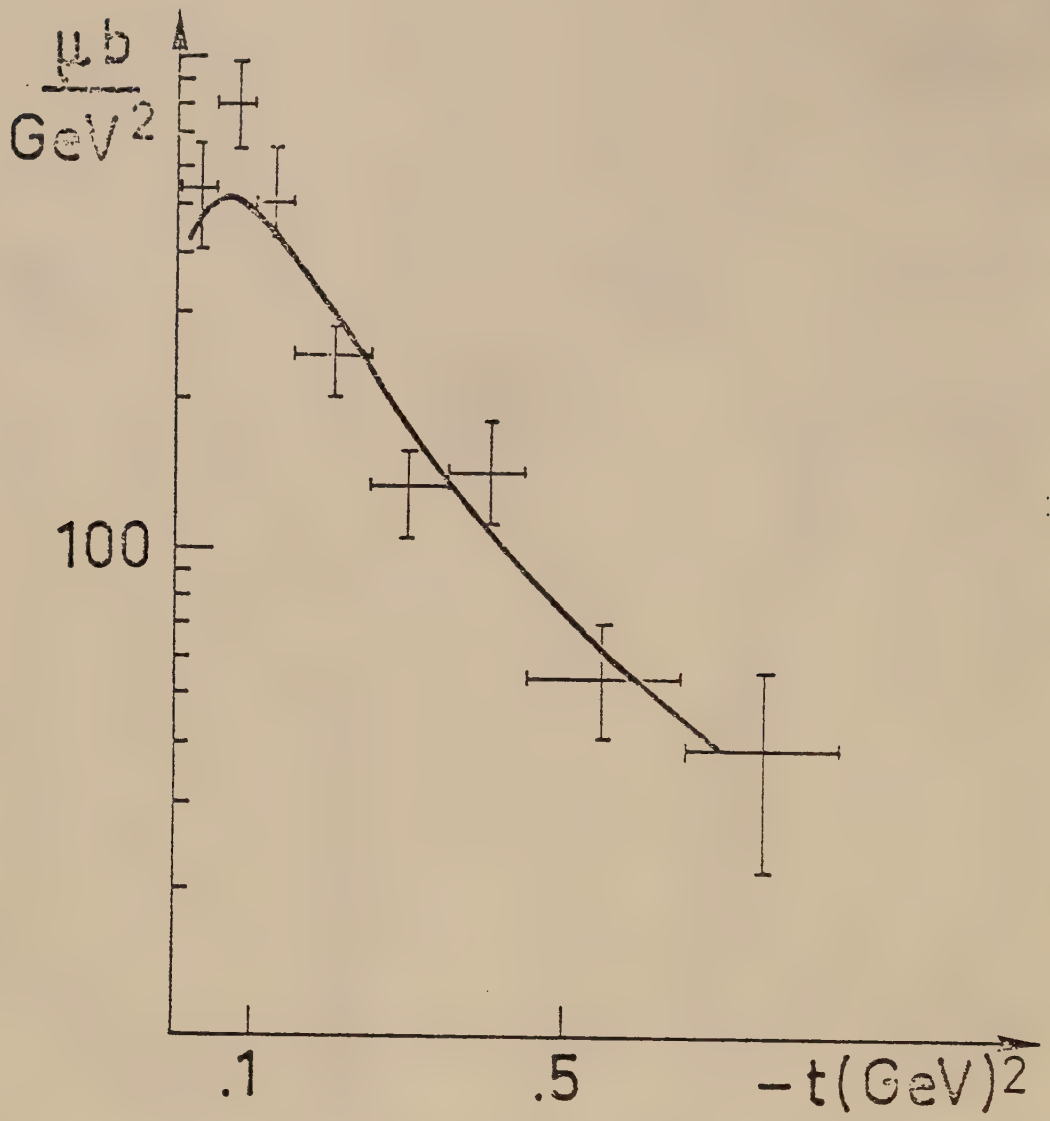


fig. 4 a

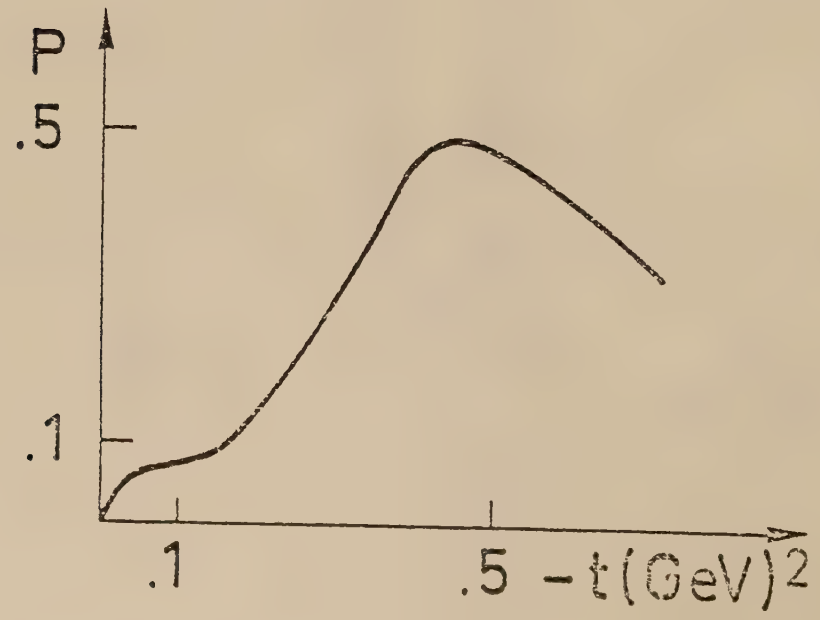


fig. 4 b

CHAPITRE V

THE EVEN SIGNATURED, ISOSCALAR EXCHANGE AMPLITUDES AT 6 GeV/c

by

G. GIRARDI, C. GODRECHE and H. NAVELET

ABSTRACT

An analysis of experimental data on $\bar{K}N$ elastic reactions and regeneration experiments leads to the $I_t = 0$ even signatured amplitudes. The results are similar to those obtained through amplitude analysis of πN scattering. Spin conservation seems to be favoured rather than helicity conservation. Predictions of the spin rotation parameters $R^{\pm}$ and of the polarisation for the regeneration experiments are given.

SECTION I

INTRODUCTION

In a previous work[1] we have studied the A_2 exchange amplitudes at 6 GeV/c by using experimental data on $\pi n \to \eta n$ and KN charge exchange scattering. Here we present a similar analysis on KN elastic reactions and regeneration experiments in order to obtain the $I_t = 0$ even-signatured amplitudes (with our conventions : $\mathbb{P}_K - \frac{1}{2} f_o$)(*) which contribute to these reactions.

In this problem we have to determine at each t-value 4 complex amplitudes $[(\mathbb{P}_K - \frac{1}{2} f_o)_\pm$ and $\omega_\pm]$ with only 5 known quantities : 3 differential cross sections and 2 polarisations. Thus we have to introduce more information by means of physical assumptions on the amplitudes. The way out we have chosen here is to assume some similar features between the ω amplitudes and the ρ amplitudes which are already known.

Within these assumptions, we have obtained the $(\mathbb{P}_K - \frac{1}{2} f_o)$ exchange amplitudes, which look like those of πN scattering[2]. We give the corresponding predictions on the spin rotation parameters R in KN elastic reactions and the polarisation for the regeneration experiment.

The paper is organised as follows : We first recall the content of our previous work and other physical informations, in section II we present the physical assumptions which enable us to unravel the problem. Section III is devoted to the determination of the amplitudes and the discussion of the results. The definitions of the amplitudes and observables may be found in an appendix.

SECTION II

WHAT IS KNOWN ?

In πN scattering at 6 GeV/c where we have a complete set of observables, several amplitudes analyses[2-5] have been done which determined the reaction amplitudes up to an overall phase. This phase ambiguity may be removed by assuming that the ρ helicity flip amplitude (ρ_-) is Regge like, a fact

---

(*) $\mathbb{P}, f_o, \omega, \rho, A_r$ refers to the exchange in the t-channel of the quantum numbers of the vacuum f_o. This is only a labelling of the allowed exchange of q.n.

which is supported by all available data and analyses. For example starting with the results of [2] and a Regge trajectory $\alpha_\rho(t) = .56\ (1+\frac{4}{3}t)$ one can stress the following features :

i) The $I_t = 0$ amplitude (P_π and f_o exchanges) is mainly non helicity flip. The non flip amplitude dominated by the Pomeron exchange is, as expected, almost pure imaginary. The relative phase of the flip and the non flip contributions is close to 180°.

ii) The $I_t = 1$ flip amplitude (ρ_-) is Regge-like whereas $Im\rho_+$ is peripheral (cross over zero around $t = -.2\ GeV^2$).

Using SU(3) relations except for the vacuum exchange, we have extracted the A_2 exchange amplitudes from the data on $\pi^- p \to \eta n$ and $K^- p \to \bar{K}^o n$ [1]. Thus, in this paper we will consider as already known quantities the ρ and A_2 contributions in KN scattering.

The available data are the following :

- the angular distributions of $K^\pm p \to K^\pm p$ [6] and $K_{o_L} p \to K_{o_S} p$ [7],
- the polarisations of elastic reactions [8].

To go further we have to add some informations to reduce the number of unknown quantities in the problem. Here we shall assume that the vector exchange (ω and ρ) amplitudes possess the same t structure and only differ by a multiplicative factor, namely

$$\omega_\pm = \gamma_\pm \rho_\pm \quad , \tag{1}$$

where $\gamma_\pm$ are real and t independent, $\rho_\pm$ are the amplitudes given by amplitude analyses [2]. This crude assumption is based on the following considerations.

i) The quantities

$$\Delta\sigma_\pi = \frac{d\sigma(\pi^- p) - d\sigma(\pi^+ p)}{\sqrt{8[d\sigma(\pi^+ p) + d\sigma(\pi^- p)]}} \quad \text{and} \quad \Delta\sigma_K = \frac{d\sigma(K^- p) - d\sigma(K^+ p)}{\sqrt{8[d\sigma(K^+ p) + d\sigma(K^- p)}}$$

which give a good estimate of, respectively, $Im\rho_+$ and $Im\ \frac{\rho_+ + \omega_+}{2}$ have been carefully studied in [9] and [10] . Both these quantities were found to behave

as Bessel functions $J_o(R\sqrt{-t})$ with the same radius, this corresponds to the same cross-over zero at $t \sim -.2$ GeV^2 and the same t dependence (see in particular fig.3 of Ref.[10]). The ratio $\Delta\sigma_K/\Delta\sigma_\pi$ yields $\gamma_+ \sim 5$.

ii) In the regeneration experiment, the phase at $t = 0$ of $\frac{1}{2}(\rho-\omega)_+$ was measured to be $-133° \pm 8°$ [11]. Within our assumption this phase would be $(\varphi_{\rho_+} - \pi) = -145° \pm 10°$ which is prettily consistent with the above value. In figure 1a we give our ω_+ amplitude, which is quite similar to the one obtained by amplitude analysis at 10 GeV/c[12].

iii) Concerning the helicity flip amplitudes there are strong evidences that absorption effects are small and do not affect the Regge behaviour (at least in the ρ and A_2 case). By extension we claim that :

$$\omega_- = \gamma_- \rho_- \quad , \tag{2}$$

since ω and ρ are vector mesons with the same trajectory.

Besides, in $\pi N \to \rho N$ reaction, one can isolate the isoscalar exchange component ($I_t = 0$, natural parity) i.e. the ω exchange, by appropriate combinations of density matrix elements. One perfectly reproduces these data by taking a pure Regge like ω in the $n = 1$ amplitude[13].

The values of γ_+ and γ_- are determined by using the data. For γ_+, the optical theorem gives :

$$\frac{\sigma_T(K^-p) - \sigma_T(K^+p)}{\sigma_T(\pi^-p) - \sigma_T(\pi^+p)} = \frac{1}{2}\,\frac{\mathrm{Im}(\rho+\omega)_+}{\mathrm{Im}\,\rho_+} = \frac{1}{2}(1+\gamma_+) = 3.0 \pm .5 \quad , \tag{3}$$

then

$$\gamma_+ = 5. \pm 1. \tag{4}$$

value, which is consistent with $d\sigma(K_L^o p \to K_S^o(p))$ in the forward direction[7].

Another estimate by R.J.N.Philipps[14] yields for the moduli $|\omega_+| \simeq 5|\rho_+|$.

Once γ_+ is surrounded, one can have a guess on γ_- by fitting $d\sigma(K_L^o p \to K_S^o p)$ [7]. We obtain a correct agreement with the data with $\gamma_- \simeq -\frac{1}{2}$ (Fig.2a). Let us stress that our assumption on the ω amplitudes allows to reproduce the regeneration experiments from the data on πN scattering

only. By taking $\gamma_+ = +5.$ and $\gamma_- = -.5$, we can predict the polarisation parameter :

$$P(K_L^o p \to K_S^o p) \cdot d\sigma(K_L^o p \to K_S^o p) = -\frac{3}{4} P(\pi^- p \to \pi^o n)\, d\sigma(\pi^- p \to \pi^o n) \quad (5)$$

The result shown in Fig.(2b) is comparable to the one of [15] .

Now we have enough information to extract the two amplitudes $M_+^o(\xi = +1)$ and M_-^o $(\xi = +1)$.

SECTION III

DETERMINATION OF THE AMPLITUDES, RESULTS

a) Following the conventions defined in the appendix, we have :

$$M_i(K^- p \to X^- p) = P_i^K - \frac{1}{2}(A_2 - \rho)_i - \frac{1}{2}(f_o \pm \omega)_i \quad . \quad (6)$$

Using the transversity amplitudes

$$M'_\pm = \frac{M_+ \pm i M_-}{\sqrt{2}} \quad (7)$$

the set of equations to be solved reads :

$$\begin{aligned} d\sigma^\pm &= |M'_+|^2 + |M'_-|^2 \\ P^\pm\, d\sigma^\pm &= |M'_-|^2 - |M'_+|^2 \quad , \end{aligned} \quad (8)$$

where $M'_\pm$ has the same form as M_i in Eq.(6) but in terms of transversity amplitudes $(P'^K)_\pm$, $(A'_2)_\pm$, $(\rho')_\pm$ $(f'_o)_\pm$ and $(\omega')_\pm$. The solutions of (8) are

$$\left| P^{K'} - \frac{1}{2} f'_o - \frac{1}{2} A'_2 \right|^2_\pm = \frac{1}{2}\left[\frac{1}{2} d\sigma^+(1 \mp P^+) + \frac{1}{2} d\sigma^-(1 \mp P^-) \right] - \frac{1}{4} |\rho' + \omega'|^2_\pm$$

$$\varphi_\pm(P'^K - \frac{1}{2} f'_o - \frac{1}{2} A'_2) = \varphi_\pm(\rho' + \omega') + \varepsilon \operatorname{Arcos} C_\pm \quad (9)$$

where $\varepsilon = \pm 1$ due to the parity of the cosine function and :

$$C_\pm = \frac{1}{2} \frac{d\sigma^-(1 \mp P^-) - d\sigma^+(1 \mp P^+)}{\left| P^{K'} - \frac{1}{2} f'_o - \frac{1}{2} A'_2 \right|_\pm \cdot |\rho + \omega|'_\pm} \quad . \quad (10)$$

Then going back to the helicity amplitudes we are left with 4 solutions for each of them, solutions which coincide at bifurcation points whenever $C_{\pm} = \pm 1$.

b) Results (Figs.3-4)

At this stage physical considerations will be helpful to single out the relevant solutions.

- at $t = 0$ the flip amplitudes must vanish, this eliminates 2 sets of solutions

- at $t = 0$ the imaginary part of the non flip amplitude has to be positive because of the optical theorem.

With these two requirements we are left with only one satisfactory solution close to the forward diection. In addition such a solution enjoys some interesting properties :

i) The real part of the non flip amplitude $M^{o}_{+}(\xi = +1)$ is very small (i.e. the Pomeron exchange which dominates the f_o exchange, is mainly imaginary) and negative in agreement with the experimental values of the real parts in the forward direction.

ii) The ratio $r = \dfrac{|M^{o}_{-}(\xi = +1)|}{|M^{o}_{+}(\xi = -1)|} \sqrt{-\dfrac{s_o}{t}}$ with $s_o = 1\ \text{GeV}^2$ is small but different from zero, $r \simeq .3$, $.4$. This indicates a violation of s-channel helicity conservation.

iii) The relative phase of $M^{o}_{-}(\xi = +1)$ and $M^{o}_{+}(\xi = +1)$ is close to 180°, then the interference term $\mathrm{Im} M^{o}_{+} M^{o*}_{-}$ poorly contributes to the polarisation. This explains why usual Regge models with no helicity flip component in M^{o} can nevertheless reproduce the polarisations of $K^{\pm} p \to K^{\pm} p$.

Properties (ii) and (iii) seem to favour spin conservation(*) rather than s channel helicity conservation. Indeed spin conservation for $M^{o}(\xi = +1)$ implies

$$\frac{M^{o}_{-}(\xi = +1)}{\sin \theta_s/2} = - \frac{M^{o}_{+}(\xi = +1)}{\cos \theta_s/2} \quad . \qquad (11)$$

(*) $g = 0$ if the scattering amplitude reads

$$M(s,t) = f(s,t) + i \vec{\sigma} \hat{n} g(s,t) .$$

At $P_{Lab} = 6$ GeV/c this implies $r \sim .3$. Such a fact was already stressed in a study of the Pomeron structure in K^+p scattering[16], where spin conservation was favoured in particular at lower energy.

One may preserve these pleasant features all along the t-range by taking advantage of the ambiguity occuring at a bifurcation point. So that we pass to another solution near $t = -.4$ GeV^2.

All our results are given in Figs.(3-4) and in table 2. The uncertainty domains were computed thanks to a powerful program[17] which takes into account correlations and physical constraints.

Obviously the flip amplitude, due to its smallness is badly known, nevertheless it appears that it is definitely non zero. A better knowledge of such a quantity could be reachable by the measurement of the R parameters in $\bar{K}N$ scattering.

Now one can also give some predictions on the R spin rotation parameters. They are found to be small negative as in πN scattering. This is also in agreement with a mean value for $R(K^-p)$ measured at 40 GeV/c[18] for $.2 < |t| < .5$: $R(K^-p) = -(.09 \pm .23)$. Despite the high value of the incident momentum this comparison is relevant since the dominance of the Pomeron in such a process allows to hope for a slow energy variation (Fig.5).

SECTION IV

CONCLUSION

We can summarize the main features of our results as follows :

assuming i) SU(3) symmetry except for the Pomeron exchange

ii) the Regge behaviour of the ρ helicity flip amplitude

iii) the identity of the amplitudes for vector exchange up to a multiplicative constant,

we have been able to extract the $I_t = 0$ even-signatured amplitudes $(p^K = \frac{1}{2} f_o)_{\pm}$ at 6 GeV/c which are very similar to the ones obtained through amplitude analyses on πN scattering.

The non flip amplitude is well determined. Its real part is small and negative.

The flip amplitude is badly determined since it is small as compared to the non flip one. Its relative phase to the non flip amplitude is close to 180°. Our results favour spin conservation rather than helicity conservation.

Predictions on the spin rotation parameters $R^{\pm}$ and on the polarisation for the regeneration experiment are given.

We recall that assumption iii) is needed in order to get the ω flip and non flip amplitudes since only five experimental measurements are available. It would be possible to get rid of this assumption, either thanks to the knowledge of three more observables (for instance, the spin rotation parameters R^{+} and R^{-} and the polarisation for the regeneration experiment), or by extracting information on the ω exchange by using data on other reactions ($\pi N \to \rho N$, $\gamma p \to \pi^{o} p$, ...).

APPENDIX

Notations and Conventions

$$\frac{d\sigma}{dt} \equiv d\sigma = |M_+|^2 + |M_-|^2$$

$$Pd\sigma = -2 \, Im \, M_+ M_-^*$$

$$Rd\sigma = -(|M_+|^2 - |M_-|^2) \cos\theta_p - 2 \, Re \, (M_+ M_-^*) \sin\theta_p$$

θ_p is the angle between the projectile and the recoil nucleon in the laboratory reference frame.

We denote $M_i(\pi^{\mp} p \to \pi^{\mp} p) = (P_\pi + f_o)_i \mp \rho_i = M_i^o \mp M_i^1$

$$M_i(K^{\mp} p \to K^{\mp} p) = (P_K)_i - \frac{1}{2}(A_2 \mp \rho)_i - \frac{1}{2}(f_o \mp \omega)_i$$

where $i = (\mp)$ refers to the non flip and the flip amplitudes. With these conventions, we have

$$M_i(K_L^o p \to K_S^o p) = \frac{1}{2}(\rho - \omega)_i \ .$$

We shall sometimes use the following notation

$$M_i^o(\xi = +1) = (P_K)_i - \frac{1}{2}(f_o)_i$$

$\xi = +1$ is the (s-u) signature.

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TABLE 1

$-t$ $(GeV)^2$	$d\sigma(K^+p \to K^+p)$ $\frac{mb}{(GeV)^2}$	$P(K^+p \to K^+p)$	$d\sigma(K^-p \to K^-p)$ $\frac{mb}{(GeV)^2}$	$P(K^-p \to K^-p)$
.04	15.7 ± 1.5	$.12 \pm .05$	$21. \pm 2.$	$.03 + .03$
.12	$10. \pm 1.$	$.18 \pm .05$	$10.5 \pm 1.$	$.07 \pm .04$
.2	$6. \pm .06$	$.21 \pm .05$	$5.6 \pm .6$	$.12 \pm .05$
.3	$3.1 \pm .3$	$.24 \pm .05$	$2.3 \pm .2$	$.2 \pm .05$
.4	$2.2 \pm .2$	$.22 \pm .05$	$1.4 \pm .15$	$.22 \pm .05$
.5	$1.5 \pm .2$	$.18 \pm .05$	$.95 \pm .1$	$.25 \pm .08$

TABLE 2

$-t$ GeV^2	$\|M\|_+$ $mb^{1/2}.GeV^{-1}$	φ_+ degrees	$\|M\|_-$ $mb^{1/2}.GeV^{-1}$	φ_- degrees
.04	3.968 ± .147	104 ± 11.3	.291 ± .107	-84.3 ± 22.6
.12	3.160 ± .105	98.1 ± 20.4	.432 ± .199	-86.9 ± 24.5
.2	2.348 ± .095	97.3 ± 22.7	.488 ± .274	-85.7 ± 26.5
.3	1.585 ± .065	116.0 ± 25.9	.388 ± .214	-74.7 ± 35.4
.4	1.269 ± .053	130.5 ± 21.2	.319 ± .155	-66.8 ± 31.1
.5	.979 ± .079	142.8 ± 27.7	.478 ± .160	-42.8 ± 26.

FIGURE CAPTIONS

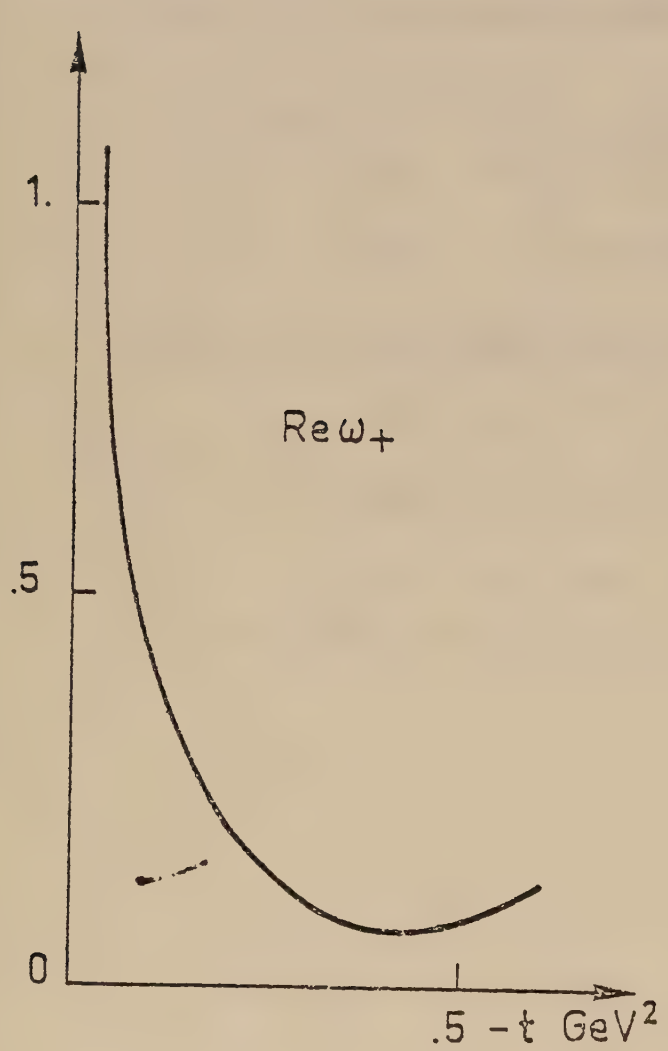
Fig.1.a-b : The non flip and flip amplitudes of the ω exchange within the assumption that $\omega_{\pm} = \gamma_{\pm}\, \rho_{\pm}$,

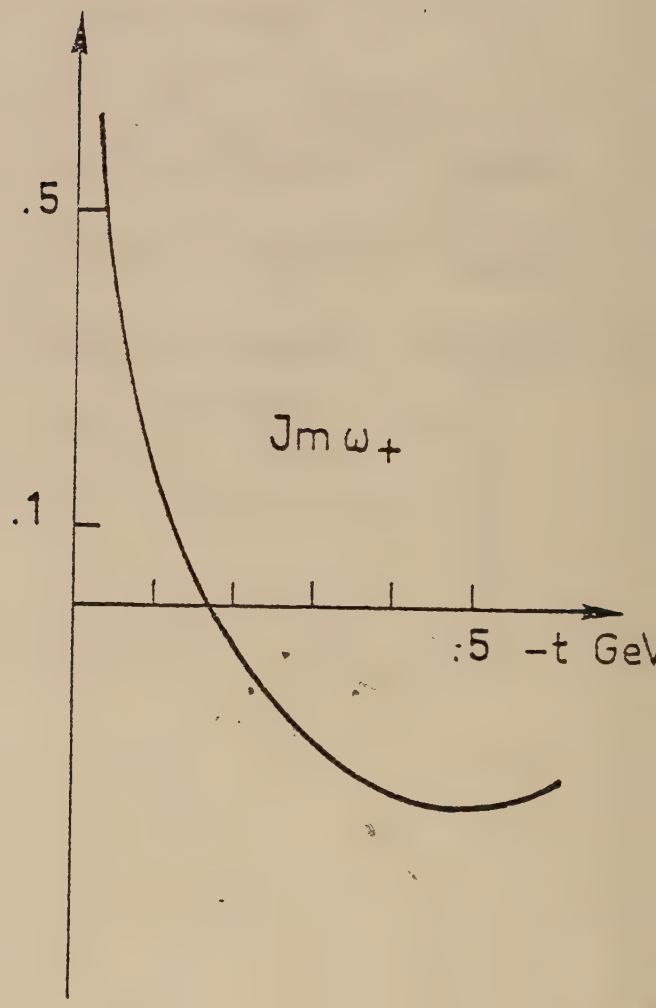
Fig.2.a-b : Angular distribution and polarisation in the regeneration experiment at 6 GeV/c . The curves correspond to $\gamma_+ = 5.$ and $\gamma_- = -.5.$

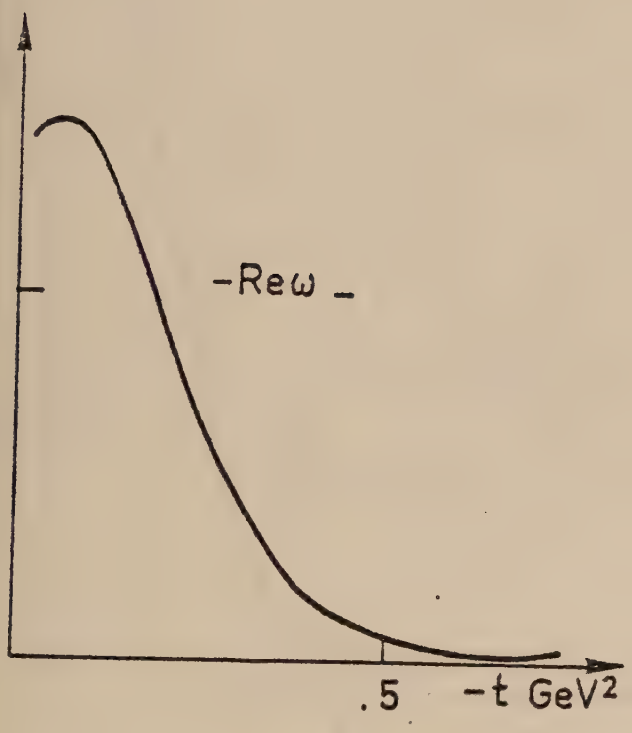
Fig.3 : Modulus and phase of M_+^o ($\xi = +1$).

Fig.4 : Modulus and phase of M_-^o ($\xi = -1$).

Fig.5 : Prediction for the spin rotation parameters of $K^{\pm}p$ elastic reactions.


1.
.5
0
Re ω+
.5 -t GeV²


.5
.1
Jm ω+
:5 -t GeV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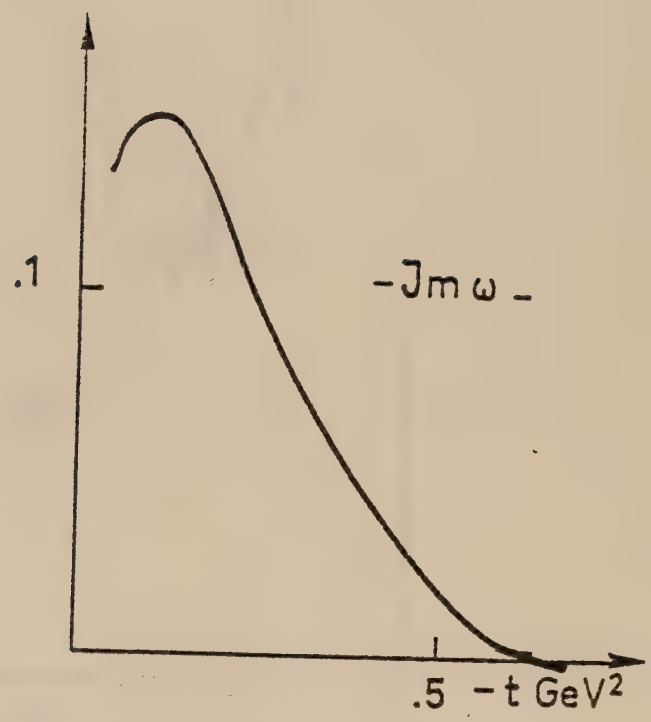
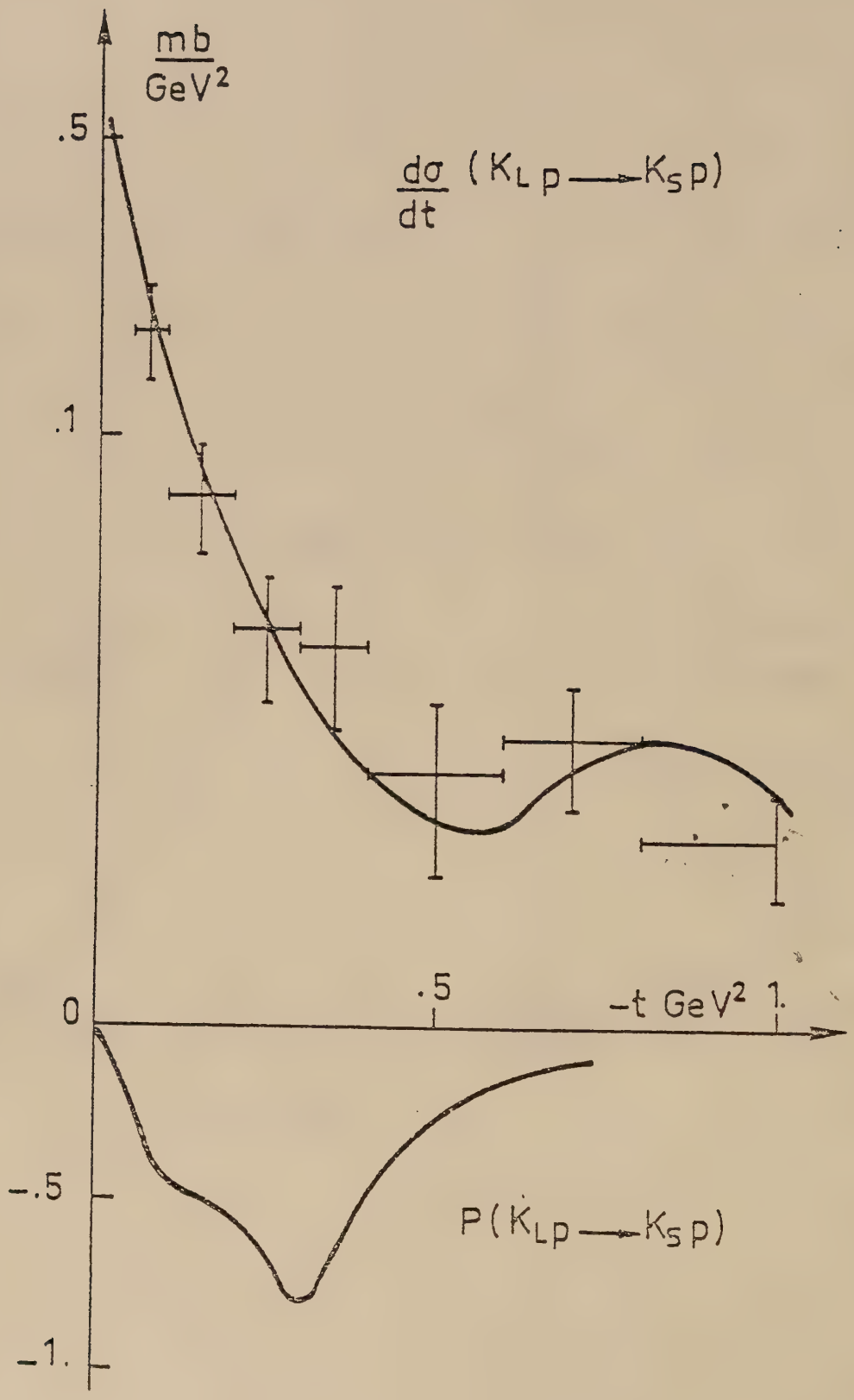
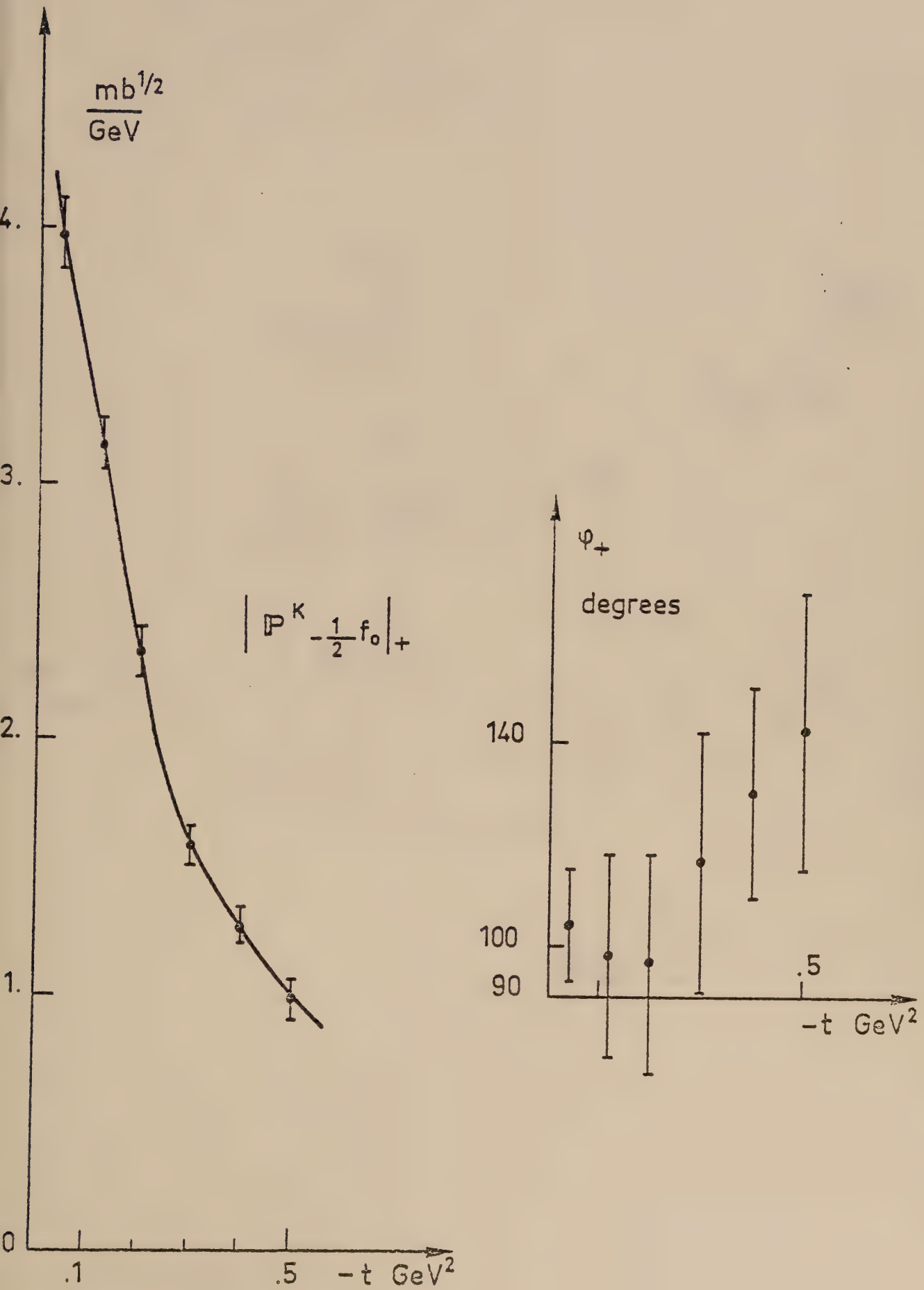
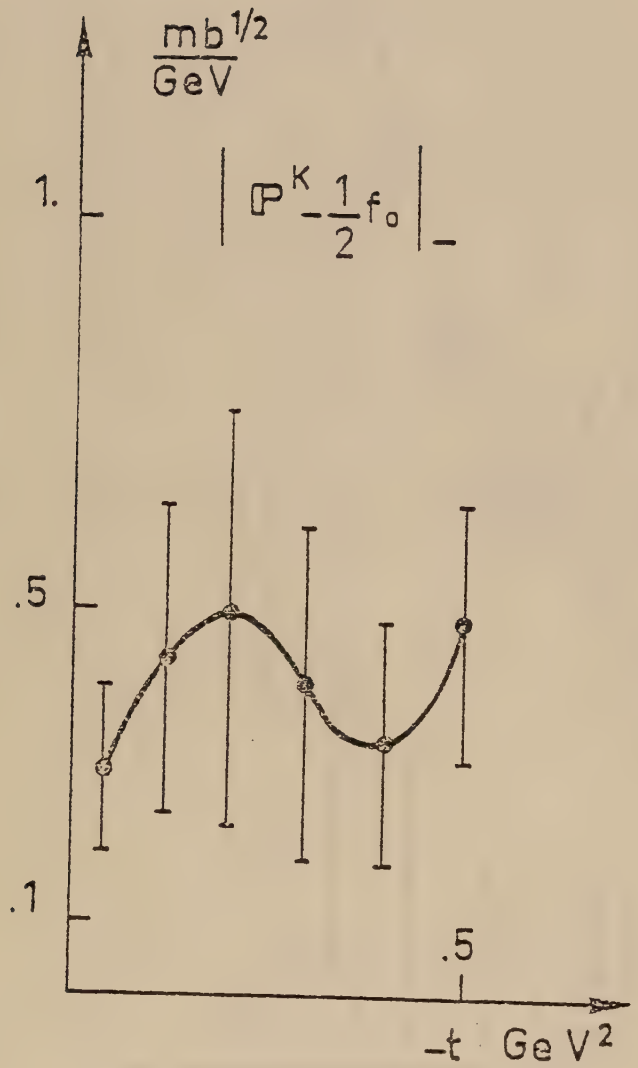
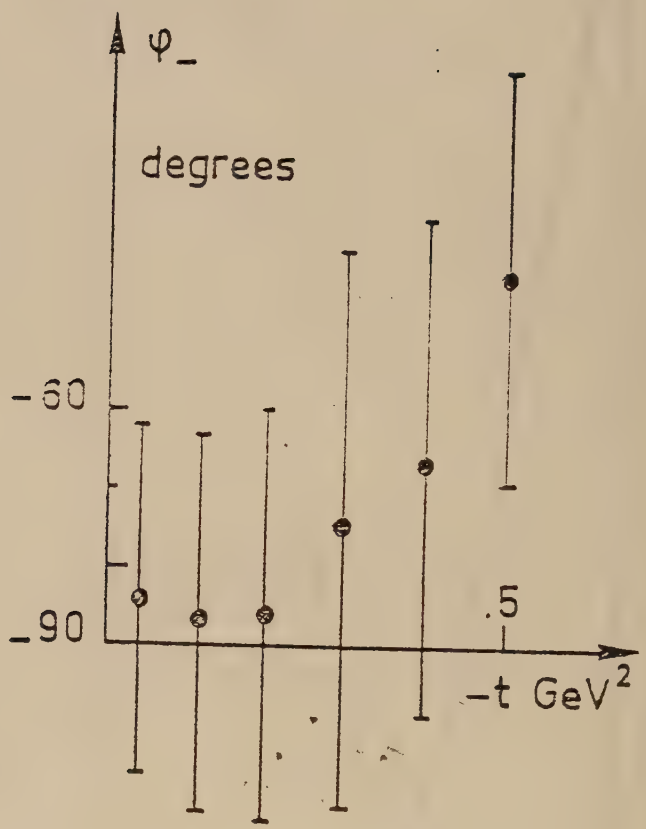
.5
-t GeV²

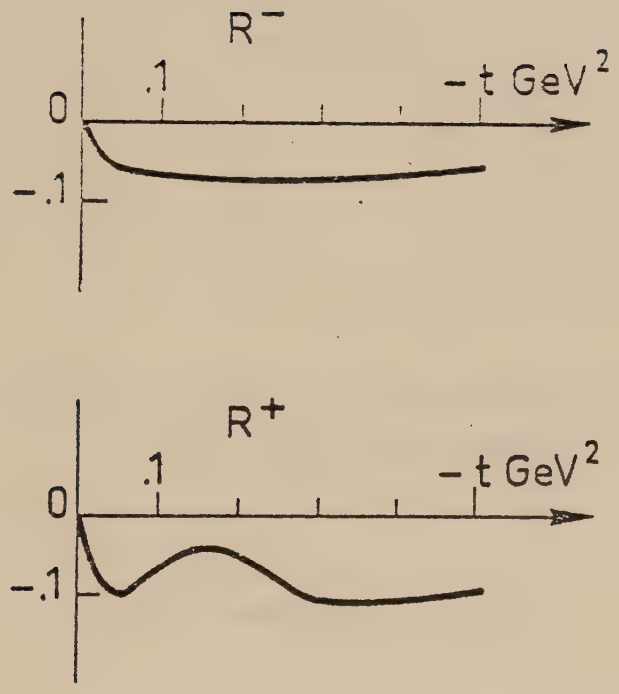

.1
-Jm ω -
.5
-t GeV²


mb
GeV²
.5
.1
0
-.5
-1.
.5
1.
-t GeV²
dσ/dt (K_L p → K_S p)
P(K_L p → K_S p)


mb^1/2
GeV
4.
3.
2.
1.
0
.1
.5
-t GeV^2
| P^K - 1/2 f_o |+
φ+
degrees
140
100
90
.5
-t GeV^2


mb^1/2
GeV
1.
.5
.1
|P^K_- 1/2 f_0|_-
.5
-t GeV^2


φ_-
degrees
-60
-90
.5
-t GeV^2


R⁻
.1
-t GeV²
0
-.1
R⁺
.1
-t GeV²
0
-.1

CHAPITRE VI

AMPLITUDE ANALYSIS AND GEOMETRICAL PICTURE FOR Δ PRODUCTION PROCESSES

by

G. GIRARDI and H. NAVELET

ABSTRACT

In this paper, a complete amplitude analysis is performed to obtain the double flip ρ and A_2 contributions. As suggested by the geometrical picture the non flip and the flip amplitudes are assumed to be identical to the ones already determined from πN and KN reactions up to t-independent real constants. Evidence for the validity of this assumption is given; SU(3) prediction works quite well with the experimental data on $KN \to K\Delta$ processes.

1 - INTRODUCTION

As far as high energy two-body interactions are concerned the Regge formalism has been a basic tool for understanding the bulk of the existing data. However to give a good description of specific experimental facts, e.g., the cross over phenomena or the polarisation in πN charge exchange, physicists have been led to create more and more sophisticated models and the original simplicity of the Regge poles model was lost. This explains the increasing interest for the amplitude analysis which acts as a substitute for the phase shift analysis at low energy where only low momentum partial waves are significant. In this approach the scattering amplitudes are directly obtained, up to an overall phase, from the experimental data if one has a complete set of observables. This is the case in $\pi N \to \pi N$ reactions at 6 GeV/c where angular distributions, polarisations and spin rotation parameters have been measured[1]. Then one can obtain the exact structure of the scattering amplitudes. For instance, the several amplitude analyses performed on πN data[2-5] tell us that the s-channel helicity flip amplitude ($n = 1$)* corresponding to t-channel isospin 1 exchange is Regge like, whereas the non helicity flip ($n = 0$) has a peripheral zero at $t \simeq .2$ GeV^2. Dual absorptive models[6] are very adequate for reproducing these features.

For KN elastic and charge exchange scattering, the experimental situation is less favourable since, for the time being, only angular distributions[7] and some polarisations[8] are available. However, it is possible by using very general assumptions such as SU(3) symmetry to obtain almost model independent amplitudes. For instance, in a previous work[9], we have obtained the amplitudes corresponding to the exchange of A_2 quantum numbers** by using SU(3) symmetry, the experimental results on η production and KN charge exchange reactions. Our results are in agreement with those of E.N. Argyres et al.[10] who obtain the ρ and A_2 exchange amplitudes with a very different approach.

In the near future, experiments using both polarised targets and beams will allow the direct determination of amplitudes in numerous quasi two-body reactions[11].

---

* n labels the net helicity flip in the s channel $n = |\lambda_3 - \lambda_4 - \lambda_1 + \lambda_2|$

** Throughout this paper, we will use the notation A_2 or ρ amplitude for amplitudes corresponding to the exchange of the A_2 or the ρ quantum numbers in the t channel.

In the present work, we are concerned with Δ production processes $\pi^+p \to \pi^o\Delta^{++}$, $\pi^+p \to \eta\,\Delta^{++}$, $K^+p \to K_o\Delta^{++}$ and $K^-n \to \bar{K}_o\,\Delta^-$ at 5.5 and 6 GeV/c. The main interest of these reactions is that, first only the ρ or the A_2 or both are exchanged and second, it is possible to analyse double flip (n = 2) amplitudes.

As afore-mentioned, one has a rather good knowledge of the ρ and A_2 (n = 0,1) amplitudes, but very little is known about the n = 2 . The aim of this work is to give information on the t-dependence of these amplitudes by using the experimental angular distributions and density matrix elements now available.

In section 2 we survey the main features of the experimental results and present the general assumptions we use to solve the problem : section 3 is devoted to the determination of the double flip ρ (resp A_2) amplitudes based on the data on $\pi^+p \to \pi^o\,\Delta^{++}$ (resp $\pi^+p \to \eta\,\Delta^{++}$). In section 4 we give the SU(3) predictions for the $K^+p \to K_o\,\Delta^{++}$ and $K^-n \to \bar{K}^o\,\Delta^-$ which agree very well with the experimental data.

2 - THE DATA AND THE PROBLEM

We first recall the notations. For any $0^-\ 1/2^+ \to 0^-\ 3/2^+$ parity conserving reactions, we have 4 independent s channel helicity amplitudes : two single flip, one non flip and one double flip

$$
\begin{aligned}
M_{\frac{1}{2},\frac{1}{2}} &= +M_{-\frac{1}{2},-\frac{1}{2}} = M_o \\
M_{\frac{3}{2},\frac{1}{2}} &= -M_{-\frac{3}{2},-\frac{1}{2}} = M_1 \\
M_{+\frac{1}{2},-\frac{1}{2}} &= -M_{-\frac{1}{2},+\frac{1}{2}} = M_1' \\
M_{\frac{3}{2},-\frac{1}{2}} &= M_{-\frac{3}{2},\frac{1}{2}} = M_2 \quad .
\end{aligned}
\tag{2.1}
$$

Where n = 0,1,2 is the net helicity flip in the s-channel. The normalisation is such that :

$$\frac{d\sigma}{dt} = |M_o|^2 + |M_1|^2 + |M_1'|^2 + |M_2|^2$$

$$\rho_{33} \frac{d\sigma}{dt} = \frac{1}{2} \{ |M_1|^2 + |M_2|^2 \}$$

$$\text{Re } \rho_{31} \frac{d\sigma}{dt} = \frac{1}{2} \text{Re } \{ M_o M_1^* + M_1' M_2^* \} \qquad (2.2)$$

$$\text{Re } \rho_{3-1} \frac{d\sigma}{dt} = \frac{1}{2} \text{Re } \{ M_o M_2^* - M_1' M_1^* \}$$

$$\rho_{33} + \rho_{11} = 1/2 \ ,$$

where $\rho_{\lambda\mu}$ labels the density matrix elements in the s-channel helicity frame.

Let us turn now to the main features of the angular distribution. It appears that the differential cross sections of reactions with a Δ in the finalstate look the same as the ones involving an outgoing nucleon.

The reaction $\pi^+ p \to \pi^o \Delta^{++}$ [12] exhibits a turn-over in the forward direction and a dip for $t \sim -.5, -.6$ GeV^2 as $\pi^- p \to \pi^o n$ does. Likewise both $\frac{d\sigma}{dt}(\pi^+ p \to \eta \Delta^{++})$ [12] and $\frac{d\sigma}{dt}(\pi^- p \to \eta n)$ present a forward turn-over and no dip near $t = -.5, -.6$ GeV^2. As in the $KN \to KN$ charge exchanges, the exotic channel $K^+ p \to K_o \Delta^{++}$ [13] lies above the non exotic one $K^- p \to \bar{K}_o \Delta^-$ [13]. The shapes of these differential cross sections are quite similar but the line reversal breaking is larger in Δ production, especially at low energy.

These facts strongly suggest that the $n = 0$ and $n = 1$, ρ and A_2 contributions in Δ production have almost the same structure as the ones occuring in πN and KN scattering [2-9], the discrepancies being due to $n = 2$ amplitudes which may give sizeable effects beyond the small t region.

This leads us to assume that the t-structure of the s-channel helicity amplitudes corresponding to well defined exchanged quantum numbers only depends on the net helicity flip $n = |\lambda_1 - \lambda_2 - \lambda_3 + \lambda_4|$. This idea has to be linked to the geometrical description [14] of high energy two body reactions. For instance Harari [14,6] has proposed that "geometry" is valid for the imaginary part of the amplitudes i.e. $\text{Im } f_n \sim J_n (R\sqrt{-t})$; in this work, we do not impose this dual absorptive model specific prescription [15]. Within this assumption

the ρ(resp.A_2) $n = 0$ and 1 amplitudes are, up to a real t-independent multiplicative factor, the same as ones which describe $\pi^- p \to \pi^o n$ (resp. $\pi^- p \to \eta n$)*. As we shall see later, this assumption is strongly supported by our results.

Now the calculation is easy to carry out. For $\pi^+ p \to \pi^o \Delta^{++}$ (resp. $\pi^+ p \to \eta \Delta^{++}$), we have to determine these 3 constants for the $n=0$ (M_o) and the two $n = 1$ amplitudes (M_1, M_1') and one complex $n = 2$ amplitude. The value of the differential cross section in the forward direction gives at once the coefficient for $n = 0$. Thus for each t' value we have to determine 4 real unknowns with the 4 experimental quantities $\frac{d\sigma}{dt}$, ρ_{33}, Re ρ_{31} and Re ρ_{3-1}.

This give us a severe test of our assumption since for each t', we have to find the same value for each coefficient of M_1 and M_1' within the experimental errors.

3 - ANALYSIS OF $\pi^+ p \to \pi^o \Delta^{++}$, $\eta \Delta^{++}$

a) The ρ exchange amplitudes

We define

$$M_o = a_o \rho_+$$

$$M_1 = a_1 \rho_-$$

$$M_1' = a_1' \rho_-$$

where the a_i's are three real constants and $\rho_\pm$ the $n = 0,1$ amplitudes of [2] where the phase ambiguity is removed by assuming that the $\rho(n=1)$ has the Regge phase with a trajectory $\alpha(t) = .56\ (1 + \frac{4}{3} t)$ [9].

The system of equations (2.2) reads

$$\begin{aligned}
\frac{d\sigma}{dt} &= a_o^2 |\rho_+|^2 + (a_1^2 + a_1'^2)|\rho_-|^2 + |M_2|^2 \\
\rho_{33} \frac{d\sigma}{dt} &= \frac{1}{2}\{a_1^2 |\rho_-|^2 + |M_2|^2\} \\
\text{Re } \rho_{31} \frac{d\sigma}{dt} &= \frac{1}{2}\{a_o a_1 \text{Re}(\rho_+ \rho_-^*) + a_1' \text{Re}(\rho_- M_2^*)\} \\
\text{Re } \rho_{3-1} \frac{d\sigma}{dt} &= \frac{1}{2}\{a_o \text{Re}(\rho_+ M_2^*) - a_1 a_1' |\rho_-|^2\} \ .
\end{aligned} \tag{3.1}$$

---

* Actually t has to be replaced by $t' = t - t_{min}$.

The determination of a_o^2 is obtained by comparing in the forward direction the two angular distributions $\frac{d\sigma}{dt}(\pi^- p \to \pi^o n) = |\rho_+|^2_{t=0}$ and $\frac{d\sigma}{dt}(\pi^+ p \to \pi^o \Delta^{++}) = a_o^2 |\rho_+|^2_{t=0}$. Thus a_o is determined up to a sign. Since the system (3.1) is invariant under the change $a_o \to -a_o$, $a_1 \to -a_1$, $a_1' \to -a_1'$ and $M_2 \to -M_2$ we will choose $a_o > 0$. Once a_o is known, we have for each t a one to one correspondance between a_1, a_1', $\mathrm{Re}M_2$ and $\mathrm{Im}M_2$, and the four observables. The numerical calculations have been performed thanks to a very powerful fit program due to M. Tournarie[15].

The results are listed in table 1 and the fitted observables in figures (1-2c). One can check that a_1 and a_1' are quite t-independent all along our t range, while the data are faithfully described. The amplitudes have then the same t-structure as in πN reactions but with different constant couplings. This constancy strongly supports our geometrical assumption. Concerning the double flip exchange $\rho(n=2)$, its contribution is small* and hence is badly determined (fig.3a-d). Indeed it is possible to neglect this amplitude, as I.J. Bloodworth did[16], and obtain a reasonable agreement with the experimental data. Anyway there is evidence for a zero in the imaginary part near $t \sim -.4$ GeV^2 (fig.3d) as obtained in[17]; such a zero cannot be accounted for by a J_2 function with radius $R \sim 1$ fm. In the real part (fig.3c) whether or not there are zeroes depends on one point at $t = -.12$ GeV^2. If it is moved downwards all structure is washed out. The ratio $|a_1/a_1'| = 1.77 \pm .23$ agrees with the values found in [16-17] and is very close to $\sqrt{3}$ as predicted in the Stodolsky-Sakurai model[18].

b) The A_2 exchange amplitudes

Now we define

$$M_o = b_o A_{2+}$$
$$M_1 = b_1 A_{2-}$$
$$M_1' = b_1' A_{2-}$$

where $A_{2\pm}$ are the helicity flip and non flip amplitudes already determined in [9]. Actually in our previous work, the error bars were overestimated. A more realistic calculation[15] enabled us to noticably reduce the uncertainties and to obtain for instance that $\mathrm{Im}\, A_{2+}$ does not vanish below $t \sim -.4$, $-.5$ GeV^2. These new results are given in table 2 and in figs.(11-12).

---

* This might be due to the kinematical t factor.

As in section III.a, b_o^2 is fixed by the differential cross sections in the forward direction, $\frac{d\sigma}{dt}(\pi^- p \to \eta n)$ and $\frac{d\sigma}{dt}(\pi^+ p \to \eta \Delta^{++})$. Since these quantities are poorly determined, b_o^2 is correspondingly uncertain. Then we have performed two different calculations for two values of b_o^2 (tables 3 and 4).

Due to the dominance of $(A_2)_-$, this does not affect b_1 and b_1', but here again these two parameters are remarquably constant in t. The corresponding observables are shown in fig.(4-5c) and the amplitudes $A_2(n=2)$ in figs.(6a-d). This contribution is not accurately known but is larger than in the ρ exchange. Indeed neglecting $A_2(n=2)$ would give a worse description than it did for the ρ analysis. We have evidence for a zero in the imaginary part (fig.6d) around $t = -.4\ GeV^2$; this was also observed in [17]. Another zero might occur at $t = -.2\ GeV^2$ but there is no clear conclusion for that. The real part seems structureless with a possible change of sign in the small t region.

Now what about exchange degeneracy between ρ and A_2 in this kind of processes ?

In [9] we found that for $n = 1$ the ρ and A_2 were almost exchange degenerate. They then contribute in the same way to $KN \to KN$ charge exchanges. Most of the line reversal breaking was accounted for by the $n = 0$ contributions. Even, more recent results on these reactions[19] agree better with our predictions.

If exchange degeneracy was almost verified, we would find that :

$$a_1 \sim \sqrt{3}\ b_1 \qquad \text{and} \qquad a_1' \sim \sqrt{3}\ b_1' \ ,$$

in fact

$$a_1 \sim 2.0\ b_1 \qquad \text{and} \qquad a_1' \sim 2.5\ b_1' \ ,$$

this indicates a violation of strong exchange degeneracy for the residues. However, this effect is not responsible for the large line reversal breaking in $(KN \to K\Delta)$ processes since the $n=1$ ρ and A_2 amplitudes have been shown[9] to be almost orthogonal.

To be more specific the size of breaking for a given n is of the following form

$$\Delta\sigma_n \sim 4 |A_{2n}| . |\rho_n| \cos [\Phi_{A_{2n}} - \Phi_{\rho_n}] \ .$$

Thus the $n = 2$ contributions which are likely out of phase at small t give the maximal $\Delta\sigma$

4 - PREDICTIONS FOR $K^+p \to K_o \Delta^{++}$ and $K^-n \to \bar{K}_o \Delta^-$ REACTIONS

Once the ρ and A_2 amplitudes have been determined, it is possible to predict all the observables for these two reactions by using SU(3) relations we recall briefly

$$\begin{aligned} M_n(\pi^+p \to \pi^o \Delta^{++}) &= \sqrt{2}\ \rho_n \\ M_n(\pi^+p \to \eta\ \Delta^{++}) &= -\sqrt{\frac{2}{3}}\ (A_2)_n \\ M_n(K^+p \to K_o \Delta^{++}) &= -\rho_n - (A_2)_n \\ M_n(K^-n \to \bar{K}_o \Delta^-) &= \rho_n - (A_2)_n \ . \end{aligned} \tag{4.1}$$

The reason why we have used SU(3) relations is that the four reactions obey quite well the following SU(3) sum rules,

$$\frac{d\sigma}{dt}(\pi^+p \to \pi^o \Delta^{++}) + 3\frac{d\sigma}{dt}(\pi^+p \to \eta\ \Delta^{++}) = \frac{d\sigma}{dt}(K^+p \to K_o \Delta^{++}) + \frac{d\sigma}{dt}(K^-n \to \bar{K}_o \Delta^-)$$

$$\rho^{\pi^o}_{ij}\frac{d\sigma}{dt}(\pi^+p \to \pi^o\Delta^{++}) + 3\,\rho^{\eta}_{ij}\frac{d\sigma}{dt}(\pi^+p \to \eta\,\Delta^{++}) = \rho^{K_o}_{ij}\frac{d\sigma}{dt}(K^+p \to K_o\Delta^{++}) + \rho^{\bar{K}_o}_{ij}\frac{d\sigma}{dt}(K^-n \to \bar{K}^o\Delta^-)$$

(4.2)

where ρ^a_{ij} is the density matrix element of the process a .

Experimentally, the differential cross section for the exotic reaction is larger than that for the non exotic, especially at low energy. This effect is even larger than the one observed in the two KN charge exchange reactions This is again the consequence of a violation of the exchange degeneracy for the ρ and A_2 amplitudes $n = 0$ and $n = 2$ (for $n = 1$ we have found[9] that the violation was small).

Using definition (4.1), we give in Figs.(7-10c) the prediction obtained by taking the mean values of a_i and b_i at face value. We reproduce quite well the data which are available at $p_{lab} = 5,6$ GeV/c namely the two differential cross sections[13] and the two sets of density matrix elements[20-23] measured in the Gottfried and Jackson frame[24]. Our amplitudes give the correct difference between the two differential cross sections at 6 GeV/C. As far as the density matrix elements are concerned the predictions are quite convincing, especially if we forget the measured quantities which lie outside the polarisation domain[25]. For the 3/2 case, this domain is defined by the following inequality

$$(\rho_{33} - \tfrac{1}{4})^2 + \mathrm{Re}\,\rho^2_{31} + \mathrm{Re}\,\rho^2_{3-1} < \frac{1}{4}\ , \tag{4.3}$$

which tells us that for instance $|\mathrm{Re}\ \rho_{31}|$ or $|\mathrm{Re}\ \rho_{3-1}|$ has to be less than 1/4 and the two cannot be simultaneously close to .2.

These results seem to give strong support for the validity of our assumption. It would have been possible to determine directly from the eight observables the $(A_2 \pm \rho)$ amplitudes without using the SU(3) relations (4.1). However the experimental accuracy is rather poor (in fact some data lie outside the permitted polarisation domain) and thus we have choosen to give predictions by using SU(3).

5 - CONCLUSION

To summarise, we have shown that our geometrical assumption is strongly supported by experimental data. Then we have determined the double flip ρ and A_2 amplitudes and we have predicted, thanks to the SU(3) symmetry, the angular distributions and the density matrix elements in the Gottfried and Jackson frame for the two $\bar{K}N \to K\Lambda$ processes. The agreement is quite convincing.

Some comments are in order:

The success of our assumption confirms the relevance of the s-channel net helicity flip for the description of amplitudes corresponding to a given set of quantum numbers exchanged in the t-channel.

Some other amplitudes analyses using different methodshave enabled the extraction of very common features for amplitudes corresponding to vector and tensor exchange in the t-channel.

i) For given n, the vector exchange amplitudes $[\rho, \omega, K^*_v]$[2-5,26-28] have the same structure, namely :

- For $n = 0$, there is a peripheral zero in the imaginary part corresponding in the D.A.M. to the first zero of $J_0(R\sqrt{-t})$.
- For $n = 1$, they are Regge like. Here we have to point out that the $\omega(n=1)$ amplitude in $\pi N \to \rho N$ is Regge like although it is coupled to a non flip $N\bar{N}$ vertex[26,9].

ii) The tensor exchange amplitudes (f_o, A_2, K^*_t)[9,10,27-22] have also the same structure for a given n, namely.

- For $n = 0$, there is evidence for a double zero in the imaginary part near $|t| \sim .45\,\mathrm{GeV}^2$ and a simple zero in the real part at the same t value

as solution 1 of Ref.[9] does (see Fig.11a-b of this paper). Furthermore the phase of the non-flip tensor amplitude is likely a Regge phase. Then a non compensation mechanism[30] might explained these features.

- The $n = 1$ amplitudes have approximatively a Regge behaviour at least in the small t region.

Both the geometrical assumption as well as the similar behaviour of the vector or tensor amplitudes should be verified in many other reactions. They might then provide a scheme for understanding two body and quasi two-body inelastic reactions. It would be interesting to have an understanding of these features on a more theoretical basis[31].

Concerning the $n = 2$ amplitudes, they do not seem to obey the D.A.M. prescription since, within the errors, the imaginary part zeros are not connected with those of the $J_2(R\sqrt{-t})$ at least for the usual value of R.

The kinematical t factor makes their determination difficult since we have only credible measurements up to $t = -.6$ GeV^2. More accurate experimental results at large t value are indeed welcome.

We recall that the geometrical assumption is needed to get the $n = 2$ amplitudes since only three density matrix elements are available. Experiments with polarised targets would allow a complete amplitude determination[11] and thus would test directly this assumption and also the validity of the SU(3) relation(4.1)/

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J.P. ADER et al. who propose a model with b-universality namely $f_n(b,s) \propto b^n f_o(b,s)$, which fairly agrees with data for the ρ, A_2, N_α and Δ exchanges (this work is under completion).

TABLE 1

$-t$ GeV^2	a_1	a_1'	$\lvert\rho\rvert$ (n=2) $mb^{1/2}/GeV$	Φ_ρ (n=2) degrees
0.04	-1.68 ± 0.23	0.93 ± 0.24	$(5.40 \pm 4.89)10^{-2}$	77. ± 115.
0.12	-1.79 ± 0.11	0.81 ± 0.18	(6.75 ± 4.01) "	38 ± 103.
0.20	-1.69 ± 0.16	0.97 ± 0.20	(4.99 ± 4.17) "	88 ± 68.
0.30	-1.65 ± 0.20	1.0 ± 0.23	(5.44 ± 3.85) "	172 ± 42.
0.40	-1.45 ± 0.21	0.92 ± 0.26	(7.85 ± 2.08) "	213 ± 31.
0.50	-1.59 ± 0.28	0.95 ± 0.32	(7.64 ± 1.96) "	235 ± 42.

with $a_o = 1.87$:

averaged values : $\langle a_1 \rangle = -1.66 \pm 0.13$

$\langle b_1 \rangle = 0.93 \pm 0.06$

TABLE 2

$-t$ GeV^2	Re A_2 $(n = 0)$ $mb^{1/2}\ GeV^{-1}$	Im A_2 $(n=0)$ "	Re A_2 $(n=1)$ "	Im A_2 $(n=1)$ "
0.04	0.190 ± 0.053	-0.131 ± 0.047	0.262 ± 0.050	-0.240 ± 0.051
0.12	0.087 ± 0.041	-0.044 ± 0.034	0.347 ± 0.035	-0.283 ± 0.040
0.20	0.045 ± 0.025	-0.022 ± 0.019	0.378 ± 0.025	-0.235 ± 0.034
0.30	0.053 ± 0.038	-0.036 ± 0.026	0.328 ± 0.030	-0.167 ± 0.046
0.40	0.033 ± 0.022	-0.031 ± 0.021	0.274 ± 0.018	-0.112 ± 0.037
0.50	-0.020 ± 0.020	-0.017 ± 0.019	0.241 ± 0.018	0.037 ± 0.066
0.60	-.048 ± .030	-.083 ± .032	.110 ± .050	.146 ± .044
0.70	-.085 ± .036	-.095 ± .035	.065 ± .035	.059 ± .035

TABLE 2' *

0.60	-0.008 ± 0.009	-0.004 ± 0.005	0.026 ± 0.063	0.197 ± 0.026
0.70	-0.012 ± 0.025	-0.029 ± 0.036	-0.002 ± 0.05	0.156 ± 0.026

* This table corresponds to solution 2.

TABLE 3 $a_o^2 = 0.45$

$-t$ GeV^2	b_1	b'_1	$\lvert A_2 \rvert$ (n = 2) $mb^{1/2}.GeV^{-1}$	Φ_{A_2} degrees
0.04	-0.687 ± 0.113	0.391 ± 0.079	0.185 ± 0.048	267. ± 25.
0.12	-0.739 ± 0.096	0.376 ± 0.078	0.078 ± 0.080	242. ± 40.
0.20	-0.782 ± 0.089	0.394 ± 0.078	0.080 ± 0.080	193. ± 83.
0.30	-0.907 ± 0.117	0.368 ± 0.084	0.120 ± 0.091	191. ± 62.
0.40	-0.948 ± 0.161	0.364 ± 0.085	0.154 ± 0.078	217. ± 25.
0.50	-0.892 ± 0.173	0.387 ± 0.095	0.156 ± 0.06	155. ± 25.

$\langle b_1 \rangle = -0.826 \pm 0.105$; $\langle b'_1 \rangle = 0.380 \pm 0.013$

TABLE 4 $a_o^2 = 0.64$

$-t$ GeV^2	b_1	b'_1	$\lvert A_2 \rvert$ (n = 2) $mb^{1/2}.GeV^{-1}$	Φ_{A_2} (n = 2) degrees
0.04	-0.673 ± 0.106	0.395 ± 0.079	0.196 ± 0.045	266. ± 22.
0.12	-0.738 ± 0.096	0.369 ± 0.078	0.085 ± 0.081	252. ± 38.
0.20	-0.738 ± 0.089	0.389 ± 0.078	0.080 ± 0.085	204. ± 56.
0.30	-0.906 ± 0.117	0.369 ± 0.084	0.121 ± 0.09	199. ± 50.
0.40	-0.946 ± 0.161	0.365 ± 0.085	0.153 ± 0.080	218. ± 25.
0.50	-0.889 ± 0.173	0.387 ± 0.095	0.158 ± 0.06	155. ± 24.

$\langle b_1 \rangle = -0.823 \pm 0.108$; $\langle b'_1 \rangle = 0.379 \pm 0.013$

FIGURE CAPTIONS

Fig.1 - Angular distribution of $\pi^+p \to \pi^o\Delta^{++}$ at 5.45 GeV/c [12]. The solid curve is the result of calculations.

Figs.2 - a.b.c - Density matrix elements of $\pi^+p \to \pi^o\Delta^{++}$ in the s-channel helicity frame at 5.45 GeV/c [12].

Figs.3 - a.b - Modulus and phase of the $\rho(n=2)$ amplitude, as obtained from $\pi^+p \to \pi^o\Delta^{++}$ data analysis.

Figs.3 - c.d - Real part and Imaginary part of the $\rho(n=2)$ amplitude.

Fig.4 - Angular distribution of $\pi^+p \to \eta\Delta^{++}$ at 5.45 GeV/c [12].

Figs.5 - a.b.c - Density matrix elements of $\pi^+p \to \eta\Delta^{++}$ [12] in the s-channel helicity frame.

Figs.6 - a.b - Modulus and phase of the $A_2(n=2)$ contribution.

Figs.6 - c.d - Real part and imaginary part of A_2 $(n=2)$.

Fig.7 - Angular distribution of $K^+p \to K^o\Delta^{++}$. Open circles are the results of [13] at 6 GeV/c. The others from [21] measured at 5 GeV/c have been reset at 6 GeV/c assuming a Regge like s-dependence with $\alpha(t) = 1/2 + t$.

Fig.8 - Angular distribution of $K^-n \to \bar{K}^o\Delta^-$ at 6 GeV/c open circles [13]. The others results [22] at 4.9 GeV/c have been corrected as in Fig.7.

Figs.9 - a.b.c - Density matrix elements of $K^+p \to K^o\Delta^{++}$ [20] open circles [21]. Arrow marked points lie outside the positivity domain.

Figs.10 - a.b.c - Density matrix elements of $K^-n \to \bar{K}^o\Delta^-$ [22] open circles [23].

Figs.11 - a.b.c.d - Real part and Imaginary part of the A_2 exchange $n = 0,1$. The error bars correspond to a more realistic calculation of the uncertainties than the one of ref.[9].

corresponds to solution 1 of Ref.[9] all along the t range.

corresponds to solution 2 of Ref.[9] for $|t| \geq .6(\text{GeV}^2)$.

Figs.12-a.b.c.d - Modulus and phase of the $n = 0$ and 1 A_2 exchange amplitudes.

corresponds to solution 1 of Ref.[9] all along the t range.

corresponds to solution 2 of Ref.[9] for $|t| \geq .6\ GeV^2$.

These figures show that :

i) For $n = 0$ the solution 1 modulus has a minimum for $.4 \lesssim |-t| \lesssim .5 GeV^2$.

ii) At the same t region, the $n = 1$ phase, which is Regge-like, vanishes.

Fig.13 - Phase of the $n = 0$, A_2 exchange amplitude for solution 1 when the α factor has been removed.

namely $\alpha > 0$ $\quad \varphi'_{A_2} = \varphi_{A_2}$

$\alpha < 0$ $\quad \varphi'_{A_2} = \varphi_{A_2} + \pi$

where α labels the A_2 trajectory which vanishes for $.4 \leq |-t| \leq .5 GeV^2$.

Figures 13 and 12.d show that within the errors the $n = 1$ and $n = 0$ phases are almost the same.

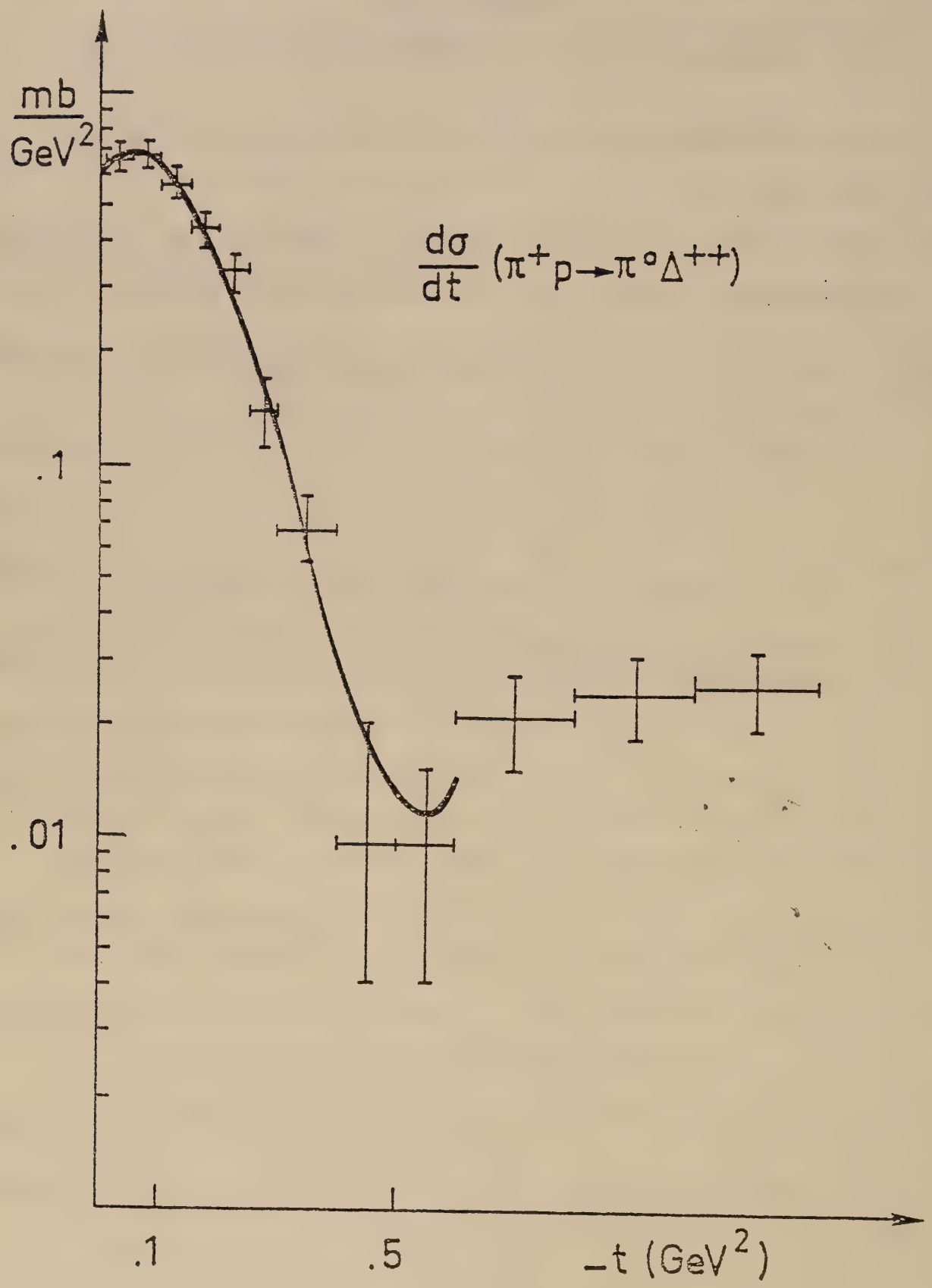


fig .1

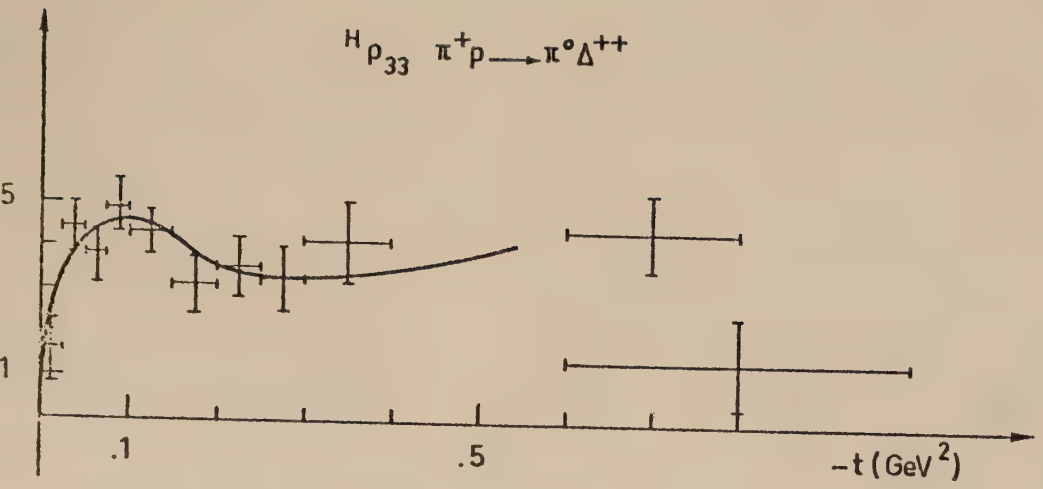


fig. 2_a

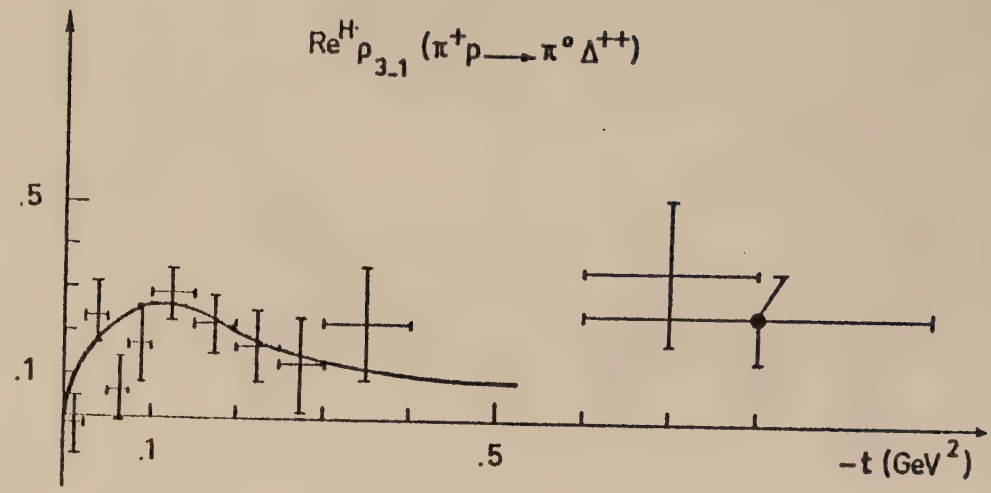


fig. 2_b

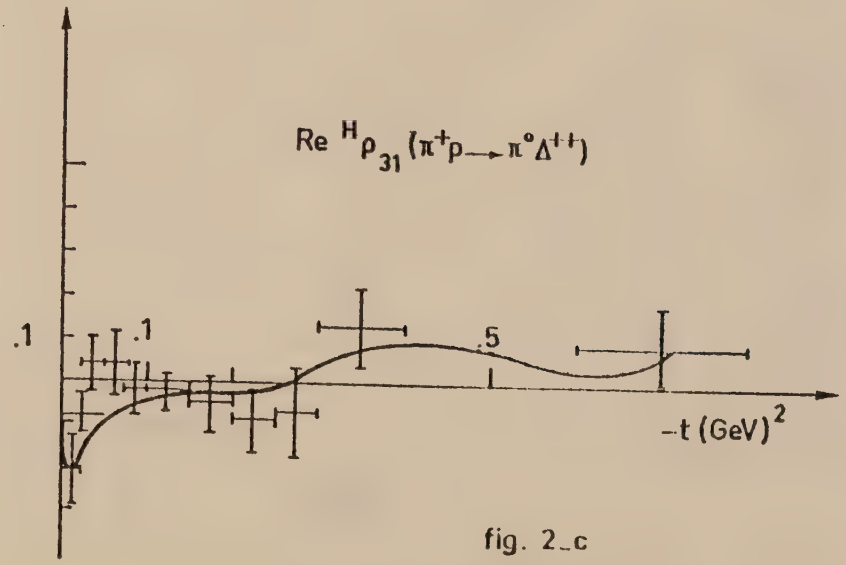


fig. 2_c

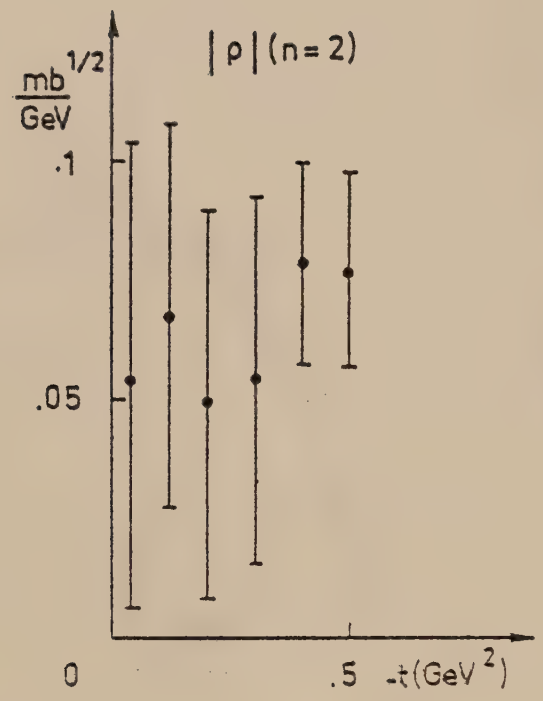


fig . 3 _a

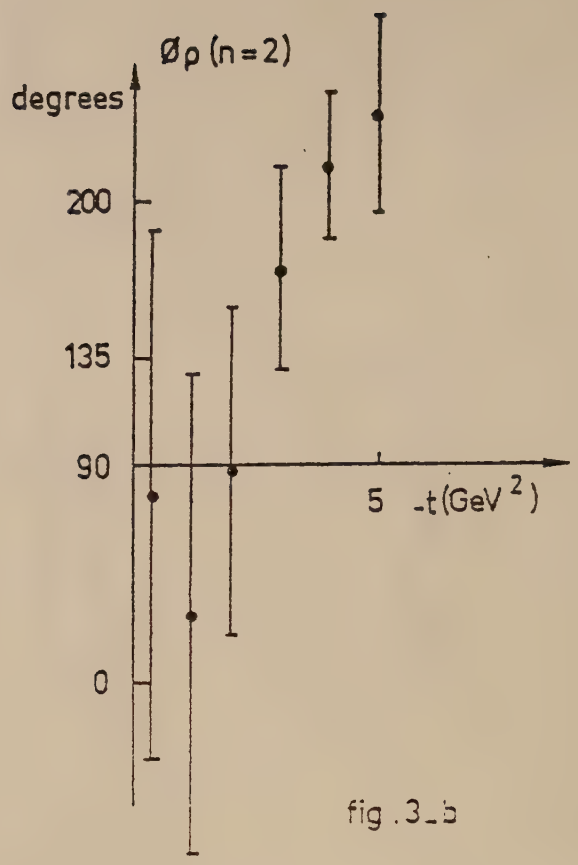


fig . 3_b

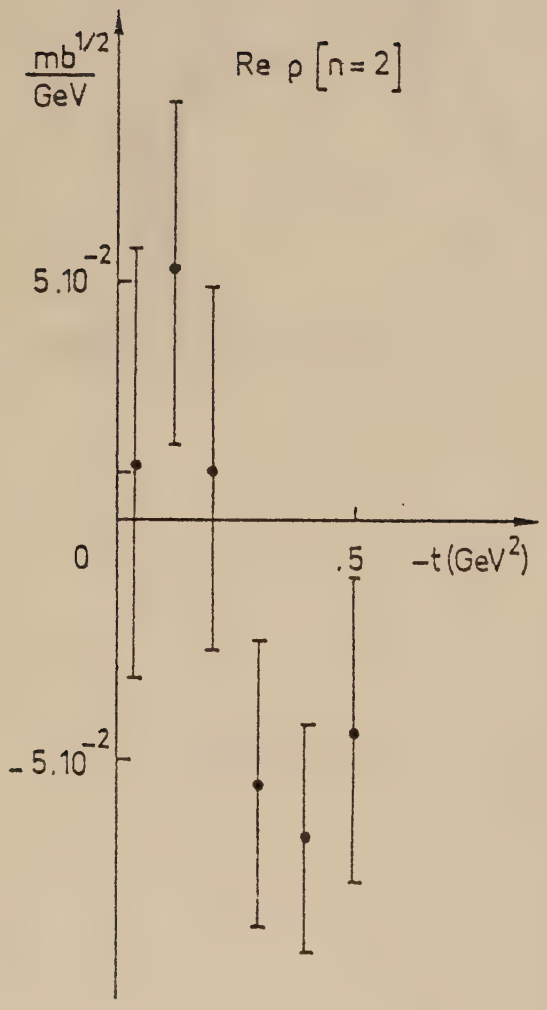


fig . 3_c

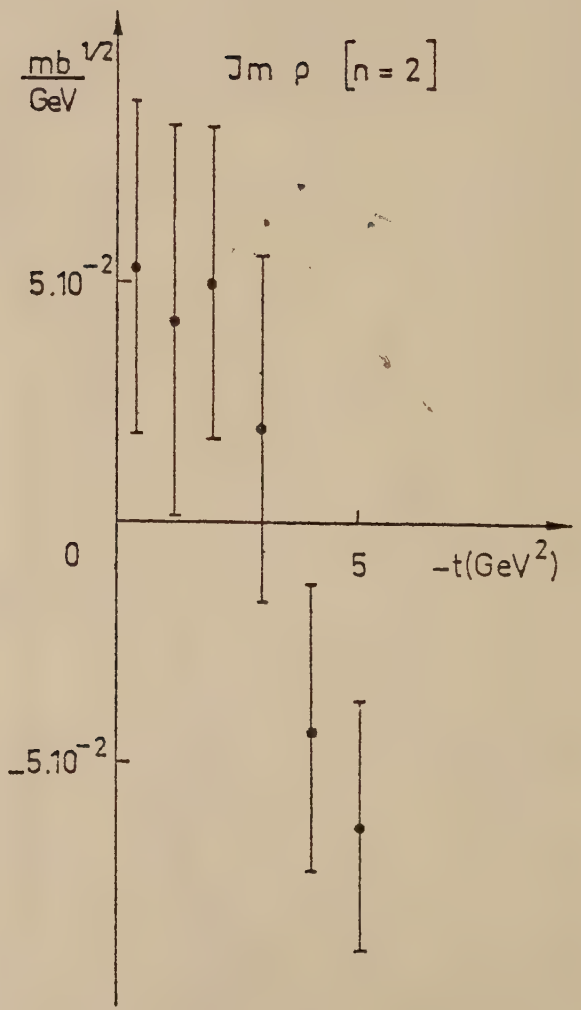


fig . 3_d

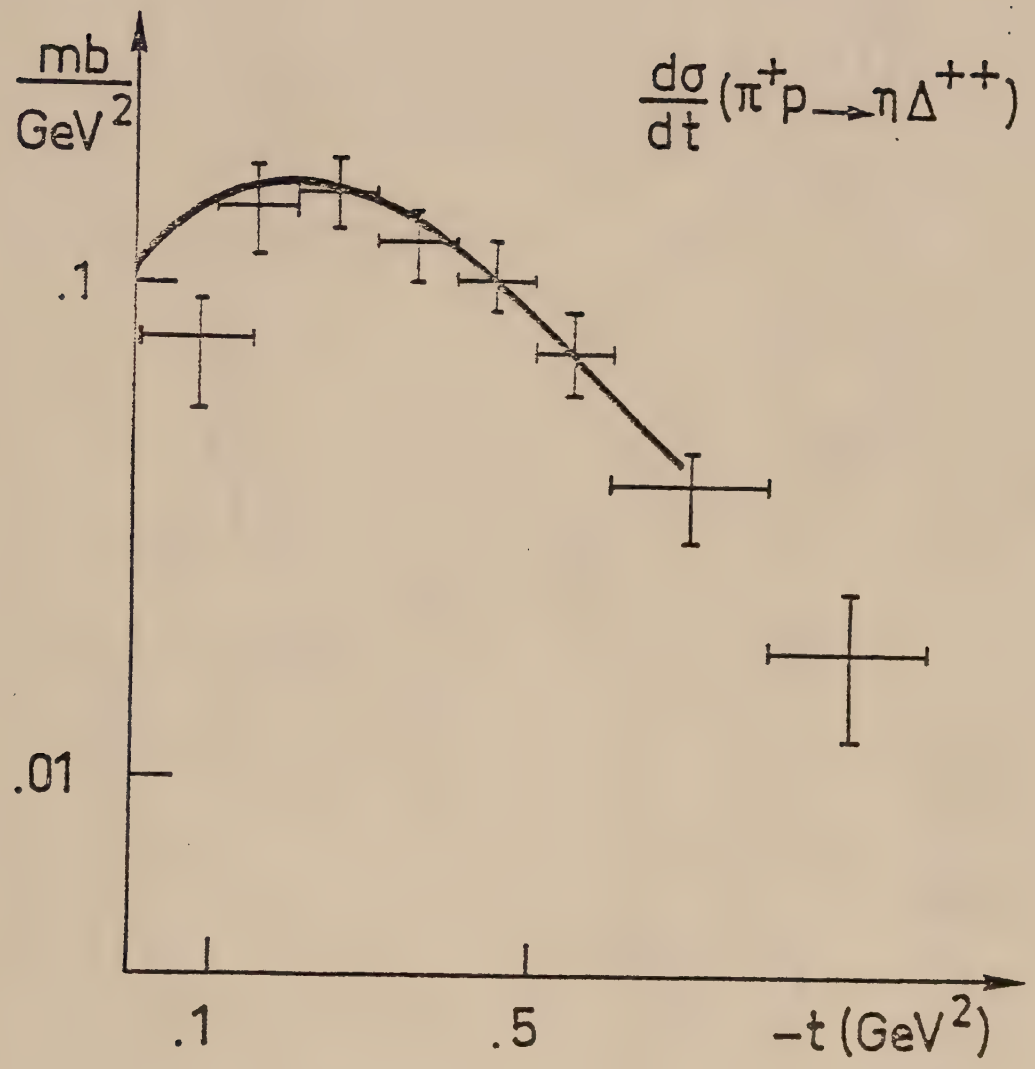


fig . 4

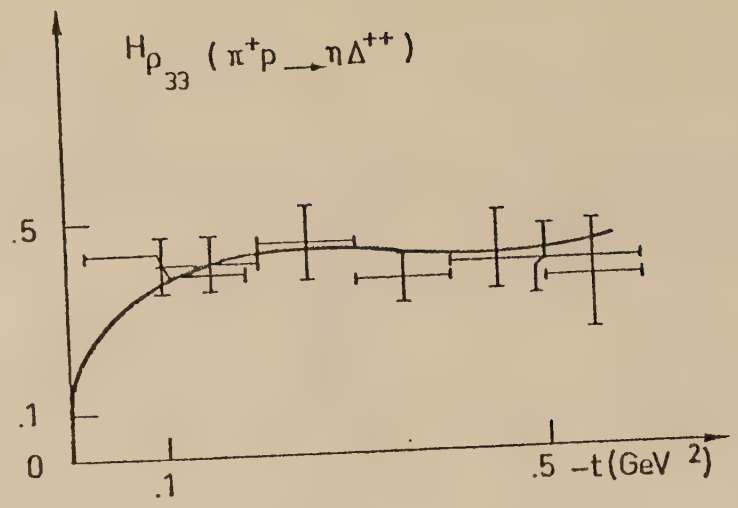


fig. 5_a

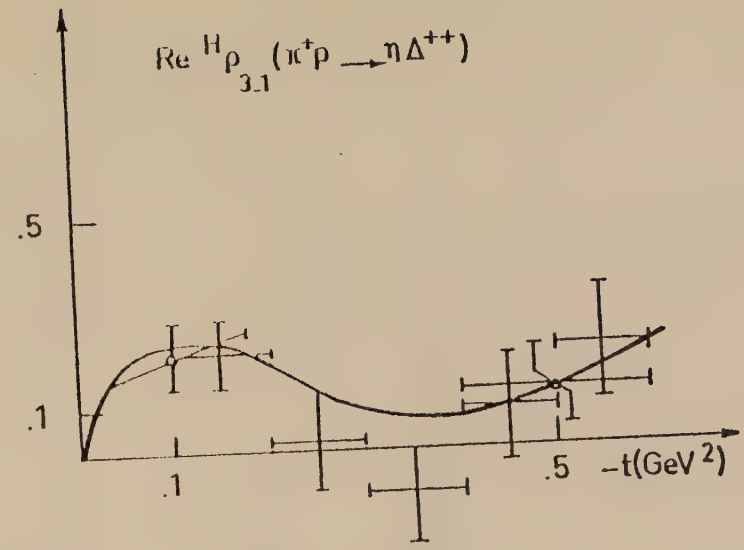


fig. 5_b

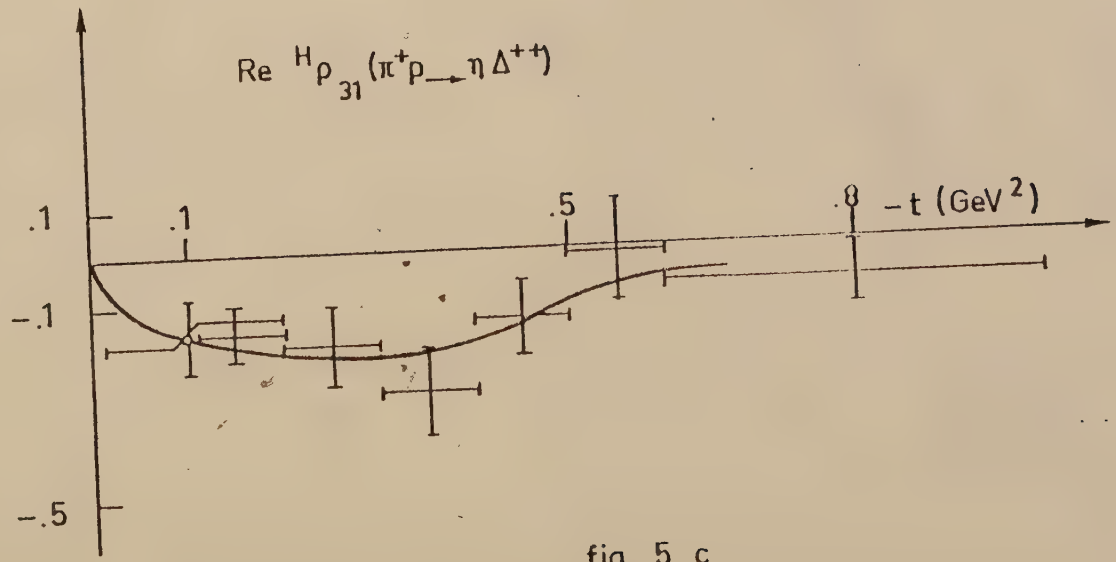


fig. 5_c

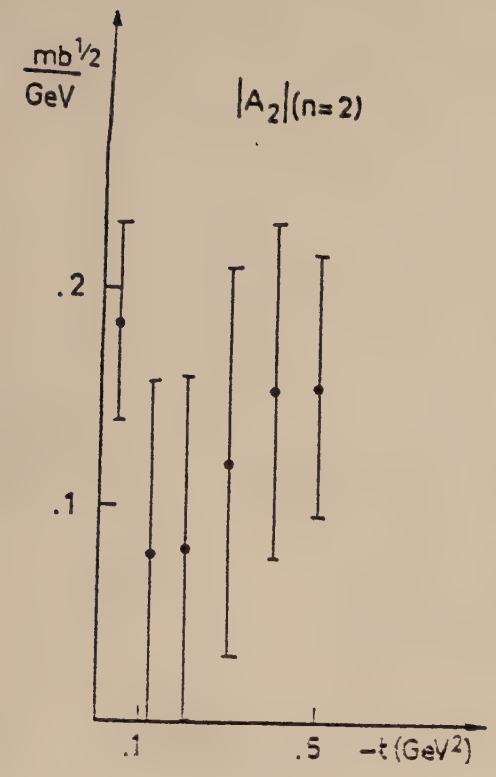


fig 6.a

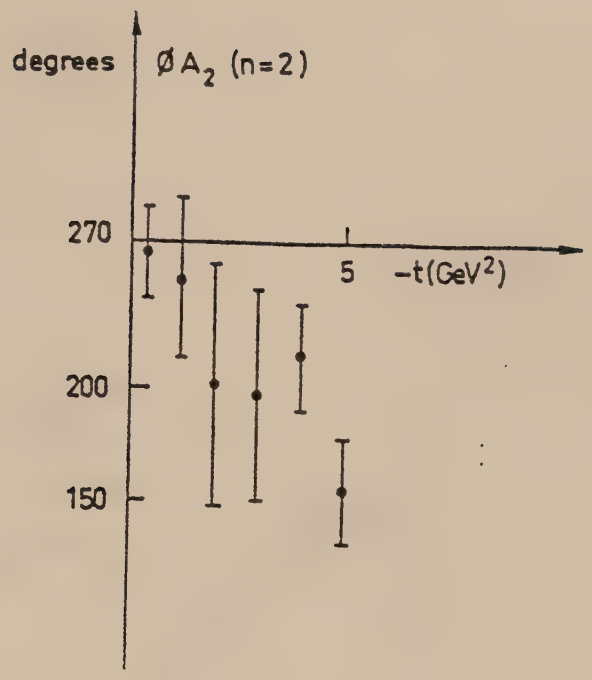


fig 6.b

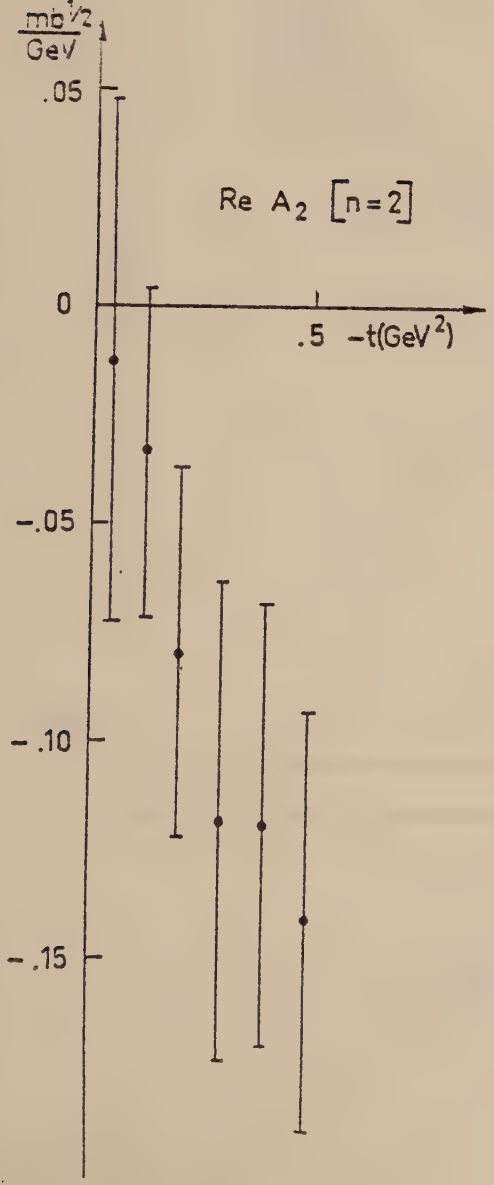


fig 6.c

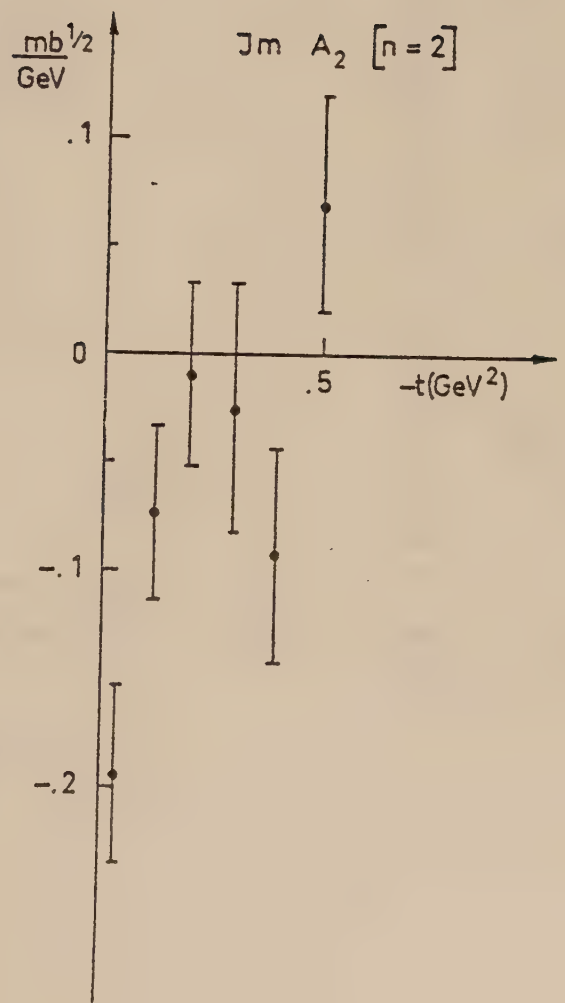


fig 6.d

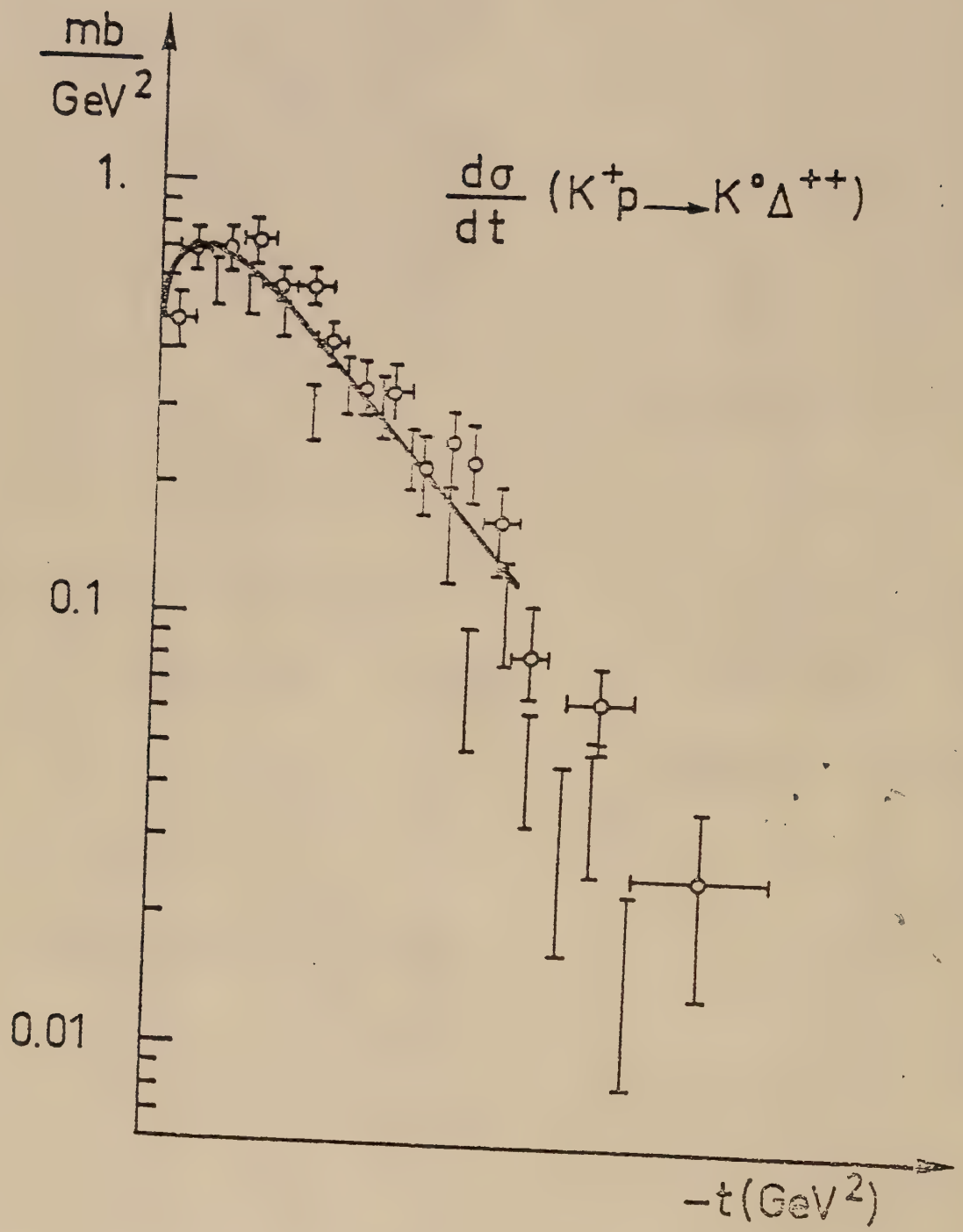


fig. 7

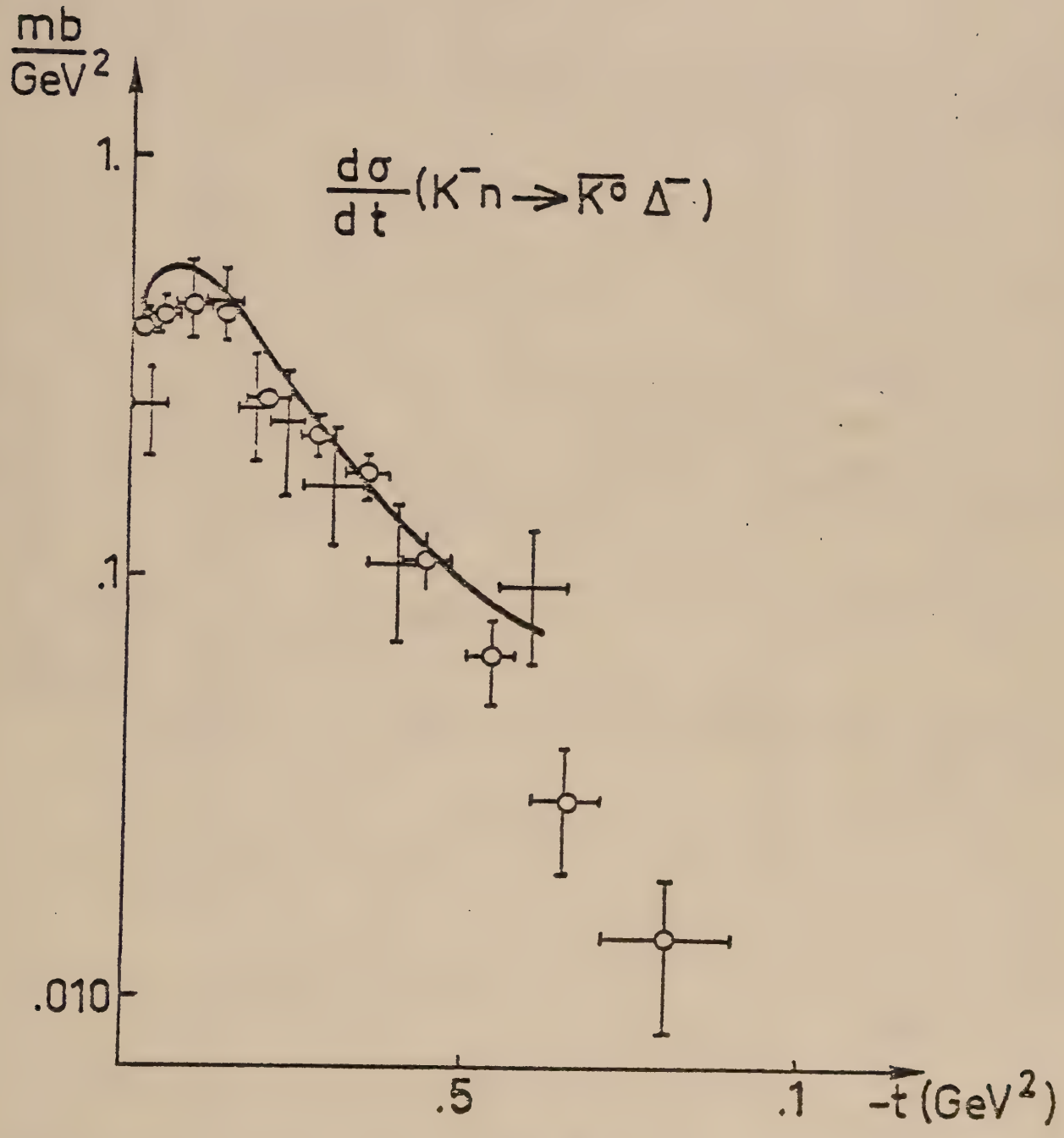


fig. 8

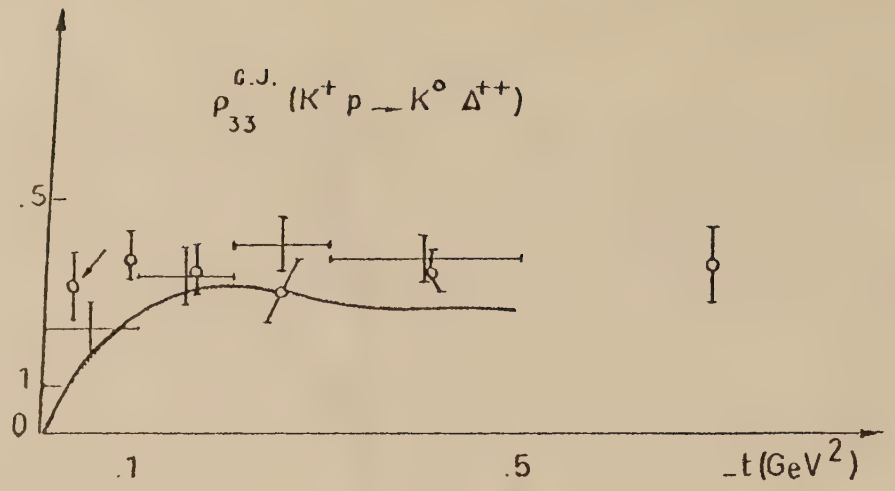


fig 9_a

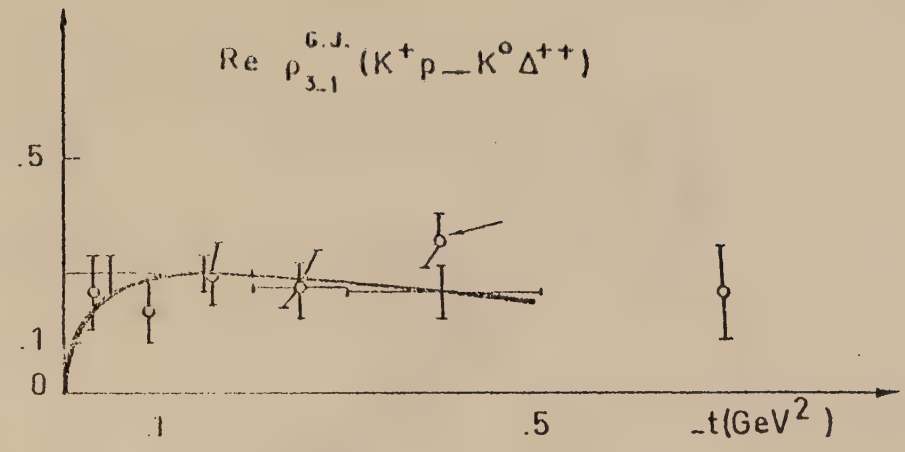


fig 9_b

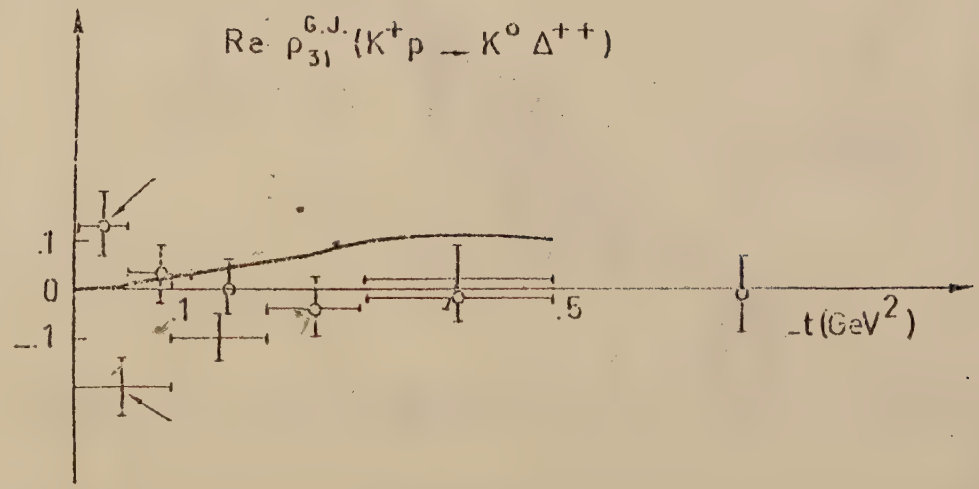


fig 9_c

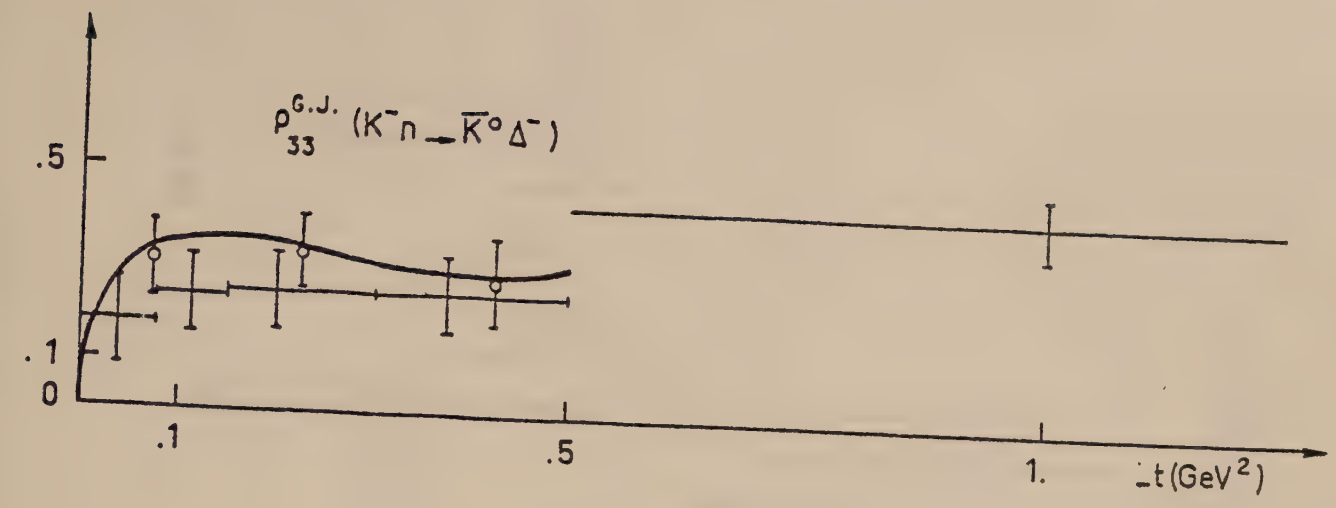


fig 10_a

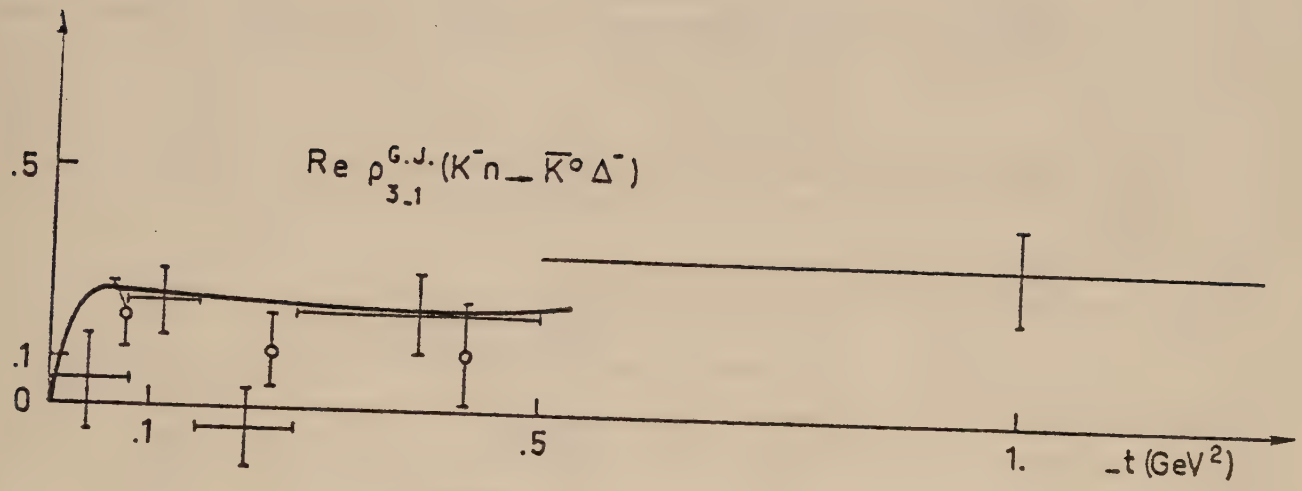
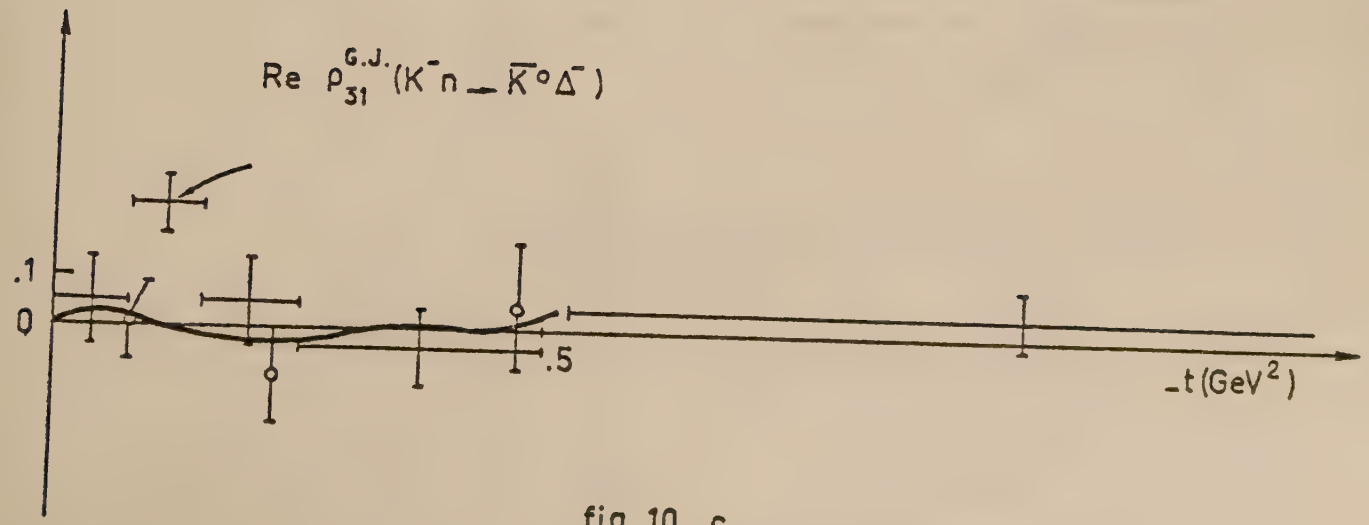


fig 10_b



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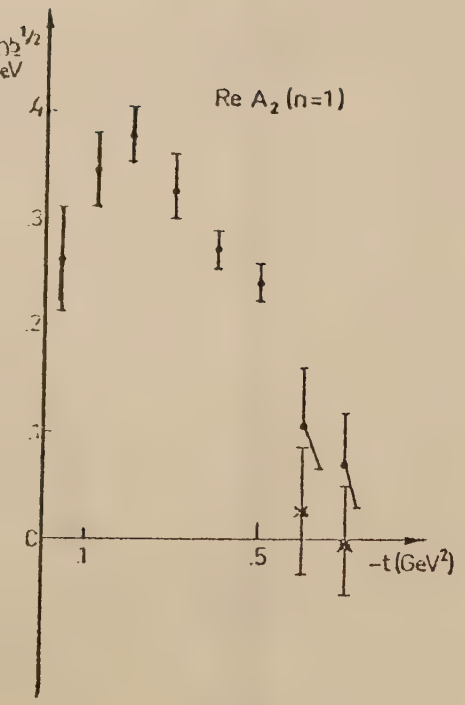


fig 11_c

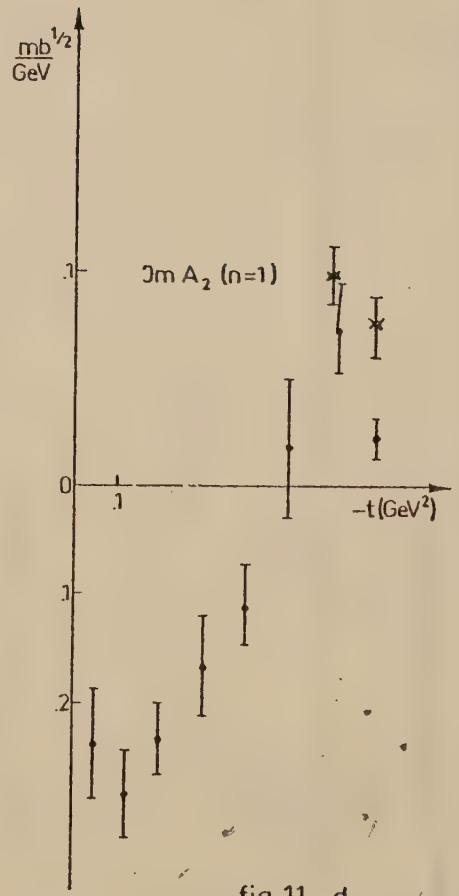


fig 11_d

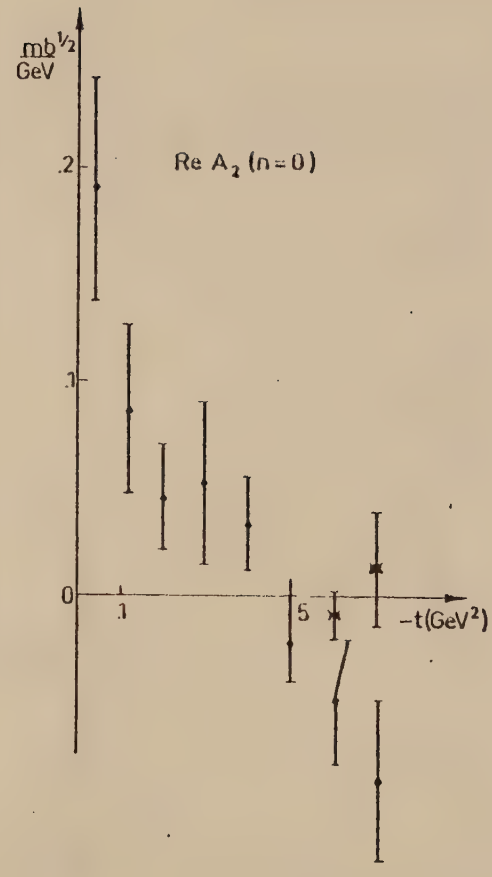
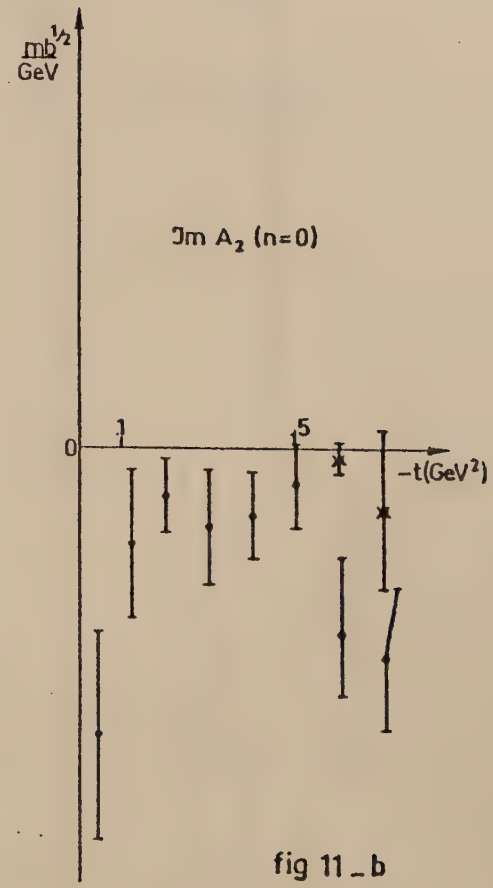


fig 11_a



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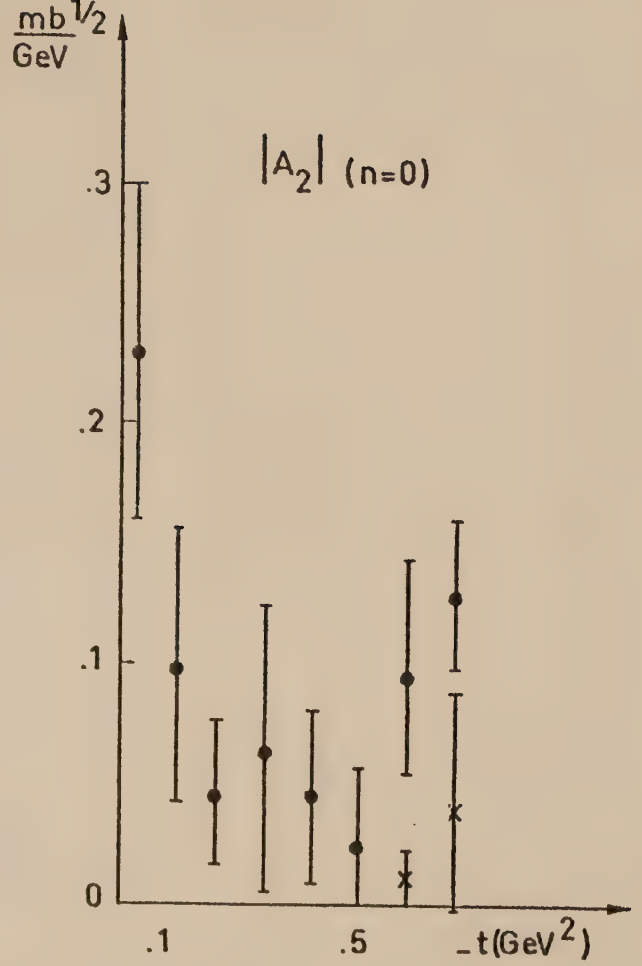

fig 12_a

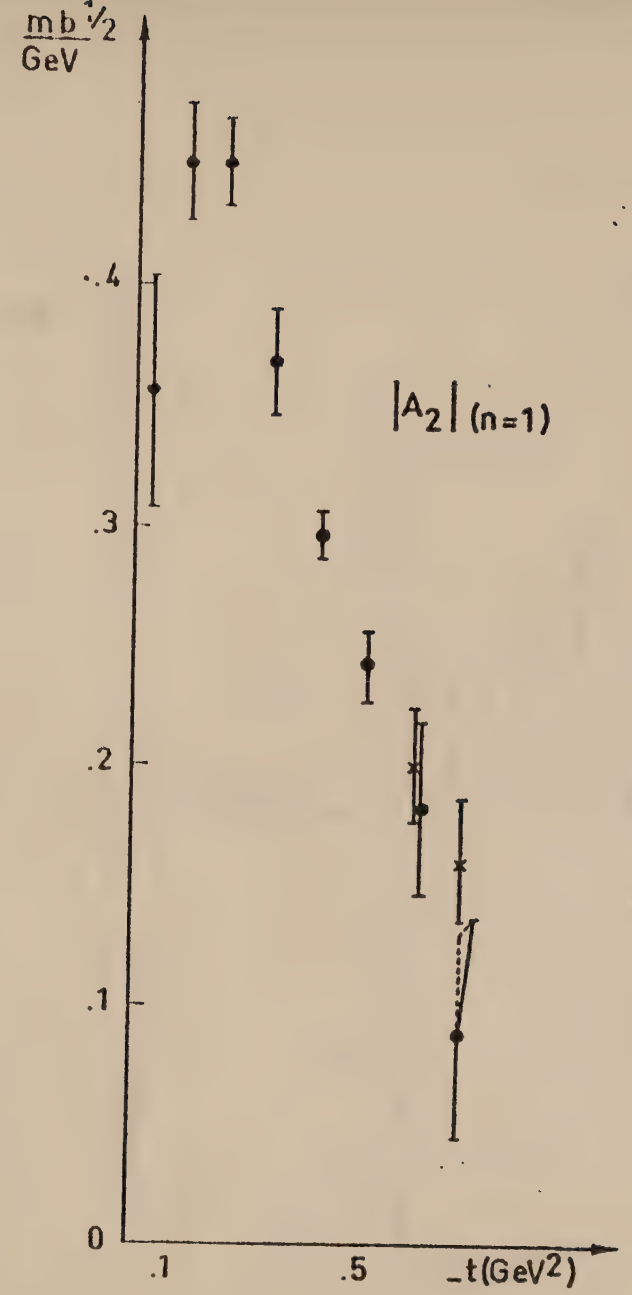

fig 12_b

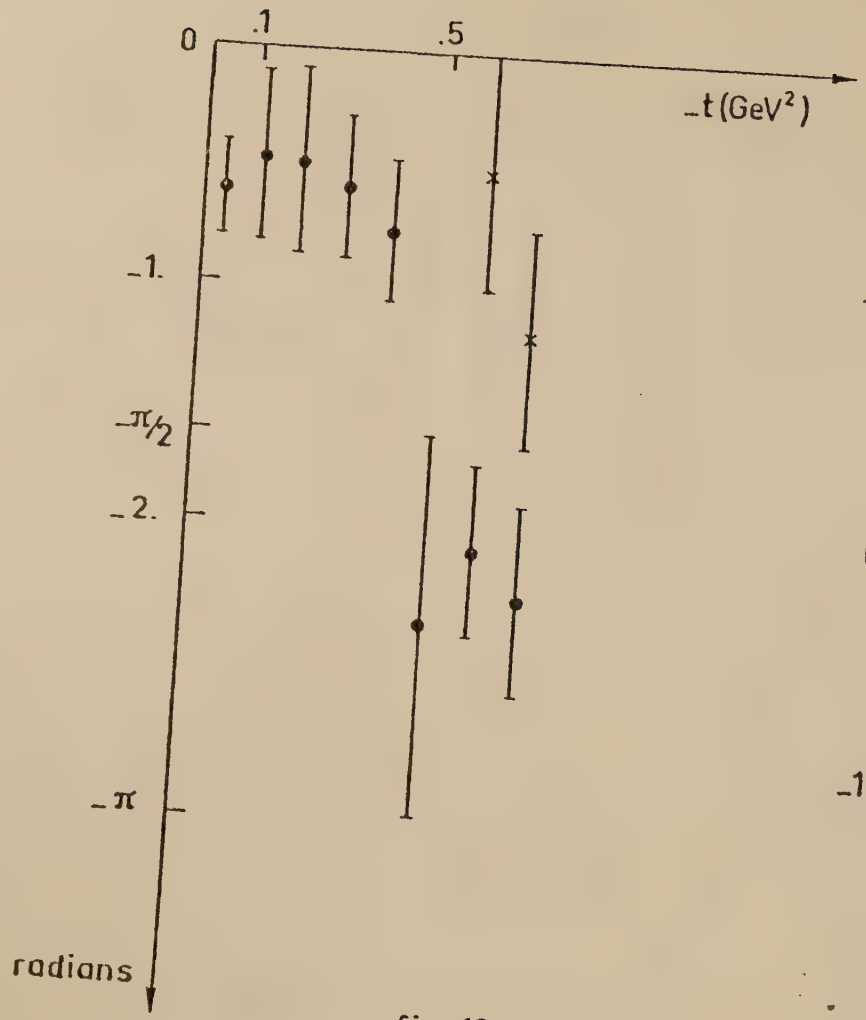

fig 12_c

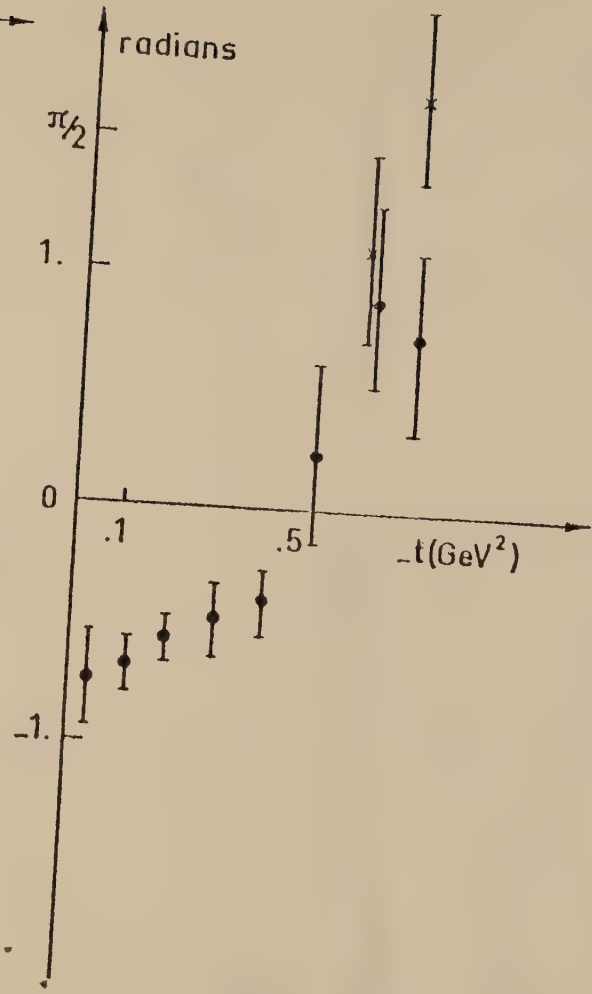

fig 12_d

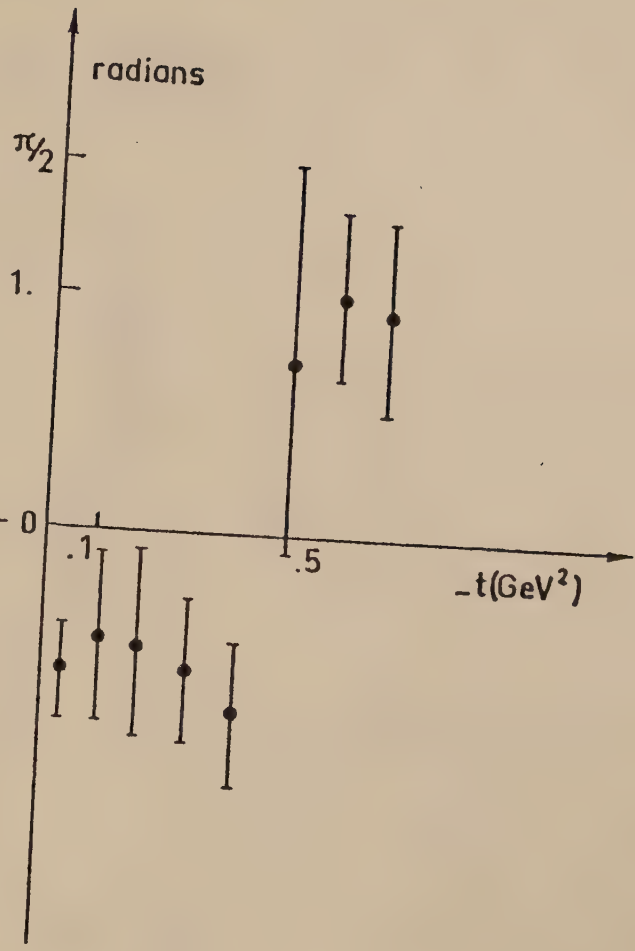

fig 13

CHAPITRE VII

A TEST OF UNIVERSAL BEHAVIOUR OF VECTOR AND TENSOR EXCHANGES

by

G. GIRARDI and H. NAVELET

A TEST OF UNIVERSAL BEHAVIOUR OF VECTOR AND TENSOR EXCHANGES

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A good description of hypercharge exchange reactions is obtained using the ρ and A_2 amplitudes analysed in πN and KN scattering, this gives a strong support to the hypothesis that vector and tensor exchange amplitudes are universal in Meson Baryon scattering.

In a recent paper [1] we presented some clues for a systematics of the vector ($J^P = 1^-$) and tensor ($J^P = 2^+$) exchange* amplitudes**. Indeed amplitude analyses [see references in 1] suggest that the t-structure of a given amplitude is determined by two elements

i) The nature of the exchange (vector or tensor).

ii) The net helicity flip in the s-channel $n = |\lambda_4 - \lambda_3 - \lambda_2 + \lambda_1|$, not with standing the possible evasive mechanisms.

This means, first, that all vector (tensor) amplitudes only differ by constant coupling factors, which are given by the SU(3) Clebsh-Gordan coefficients within a multiplet [4]. Second the t-structure will be dictated by the value of n which may correspond to several different helicity configurations of the t-channel vertices. Finally, if an amplitude is dominated by the exchange of a particle or a pure Regge pole, we shall have constraints given by residue factorisation (this is not required for more complicated exchange mechanisms).

In table I we give the expected behaviour of the various amplitudes with reference to popular and/or likely relevant models.

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*By exchange we mean : allowed set of quantum numbers in the t-channel, without any prejudice on the nature of the associated J-plane singularity.

**We use s-channel helicity amplitudes.

In this letter we shall be concerned with the hypercharge exchange reactions, i.e. with the $K_V^*(890)$ and the $K_T^*(1420)$ contributions. If $V_n(T_n)$ denote the $K_V^*(K_T^*)$ amplitudes with net helicity flip (n = 0,1) we have

$$\frac{d\sigma}{dt}(\pi N \to KY) = |T_o - V_o|^2 + |T_1 - V_1|^2$$

and (1)

$$P\frac{d\sigma}{dt} = -2\mathrm{Im}\,(T_o - V_o)(T_1 - V_1)^*$$

for the non exotic reaction ; in the line reversed process $\bar{K}N \to \pi\bar{Y}$ V_n becomes $-V_n$.

The universality assumption on vectors and tensors yields :

$$V_o = a_V \rho_o \qquad V_1 = b_V \rho_1$$
$$T_o = a_T (A_2)_o \qquad T_1 = b_T (A_2)_1 \qquad (2)$$

The ρ and A_2 contributions are those obtained in [2-3] , a and b are the adjustable couplings. A more detailed study would embody the correction arising from the difference $(\alpha_{K^*}(0) - \alpha_\rho(0))$ [4] this would affect the normalisation of the amplitude and the position of the WSNZ which is important in $K^-P \to \eta\Lambda$ reaction. However since we use experimental information only in a small energy range such a refinement is not necessary.

Equations (2) may be used as a parametrisation to fit the data, namely $\frac{d\sigma}{dt}$ and P for $\pi^+p \to K^+\Sigma^+$ (or $\pi^-p \to K^o\Sigma^o$), $\pi^-p \to K^o\Lambda$ and the line reversed reactions. The results are displayed in fig. [1.2] and show a quite good agreement. The values of a and b* are given in table II, together with the corresponding F couplings. The fluctuations of the a's and b's are very small, in agreement with our statement (i). The amplitudes in Σ and Λ reactions are related by SU(3) symmetry [4]

$$\frac{M(\pi^+p \to K^+\pi^+)}{M(\pi^-p \to K^o\Lambda)} = \sqrt{6}\,\frac{2F-1}{2F+1} \qquad \text{with} \quad F + D = 1.$$

The values $F_\pm$ are very close to those given in [4] where the authors assume that the vector and tensor exchanges have the same value of F. This is expected if we have exchange degeneracy for the K_T^* and the K_V^*. The n = 1 amplitudes cannot account for the large line reversal breaking since their phases differ by $\pi/2$ and since furthermore $(F-)_V - (F-)_T$. Concerning the n = 0 contributions, absorptive corrections by an SU(3) singlet preserve the octet relations on the couplings while giving the required breaking effect. Then the fact that $(F_+)_V$ is close to

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*Averaged over the t range.

$(F_+)_T$ is not surprising [4]. SU(3) may be used to look at the η and η' production reactions : $K^- p \to \eta\Lambda, \eta'\Lambda$ whose amplitudes are given by

$$M(K^- p \to \eta \Lambda) = \frac{1}{\sqrt{6}} (\cos\Theta + 2S_T \sin\Theta)\, T - \sqrt{\frac{3}{2}} \cos\Theta\, V$$

$$M(K^- p \to \eta' \Lambda) = \frac{1}{\sqrt{6}} (\sin\Theta - 2S_T \cos\Theta)\, T - \sqrt{\frac{3}{2}} \sin\Theta\, V$$

where Θ is the η, η' mixing angle and $S_T^2 = \frac{\sigma(\eta 8)}{\sigma(\eta 1)}$.

Since in [3] we neglected the η, η' mixing, by consistency here we shall take $\Theta = 0^\circ$. i.e. $\eta = \eta_8$ and $\eta' = \eta_1$.

Using (2) one obtains a reasonable description of the observables for these 2 reactions at least concerning the t-structure. Indeed, as already pointed out the dip in $K^- p \to \eta\Lambda$ arising from the WSNZ zero in V is too far and second the relative size of the two reactions is not correctly reproduced within our assumptions concerning the η, η' mixing*.

As a conclusion one can say that our systematics works very well for the vector and tensor exchanges and that the amplitudes used are compatible with SU(3) symmetry.

A desirable improvement would be to analyse the T exchange taking into account the η, η' mixing and to study other reactions such as K^+ photoproduction. Nevertheless it is already clear from the reactions studied here that our systematics explains quite different reactions with very few hypotheses and yields substantial information in a simple way on the underlying exchanged amplitudes.

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* For an extensive discussion, see [5].

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FIGURE CAPTIONS

Fig.1 - Hypercharge-exchange cross-sections in the region of 6 Gev/c.

a) $\frac{1}{2}(\pi^+ p \to K^+ \Sigma^+)$ 6GeV/c [6]

$\pi^- p \to K^\circ \Sigma^\circ$ 6GeV/c [7]

$K^- p \to \pi^- \Sigma^+$ 5.47 GeV/c [11]

b) $\pi^- p \to K^\circ \Lambda$ 6 GeV/c [7]

$K^\circ p \to \pi^+ \Lambda$ $5 < p_L < 8$ [14]

$K^- n \to \pi^- \Lambda^\circ$ 4.5 GeV (rescaled with $\alpha(t) = .35 + .85t$) [13]

Fig. 2 - Hypercharge exchange polarizations :

a) $\pi^+ p \to K^+ \Sigma^+$ 5.05 GeV/c [8]

7.0 GeV/c [8]

b) $K^- p \to \pi^- \Sigma^+$ 4.25 GeV/c [9]

3.95 GeV/c [10]

c) $\pi^- p \to K^\circ \Lambda$ 5 GeV/c [7]

4.5 GeV/c [12]

d) $\bar{K}^\circ p \to \pi^+ \Lambda$ 4.5, 6 GeV/c [14]

$K^- p \to \pi^\circ \Lambda$ 3.95 GeV/c [10]

TABLE I

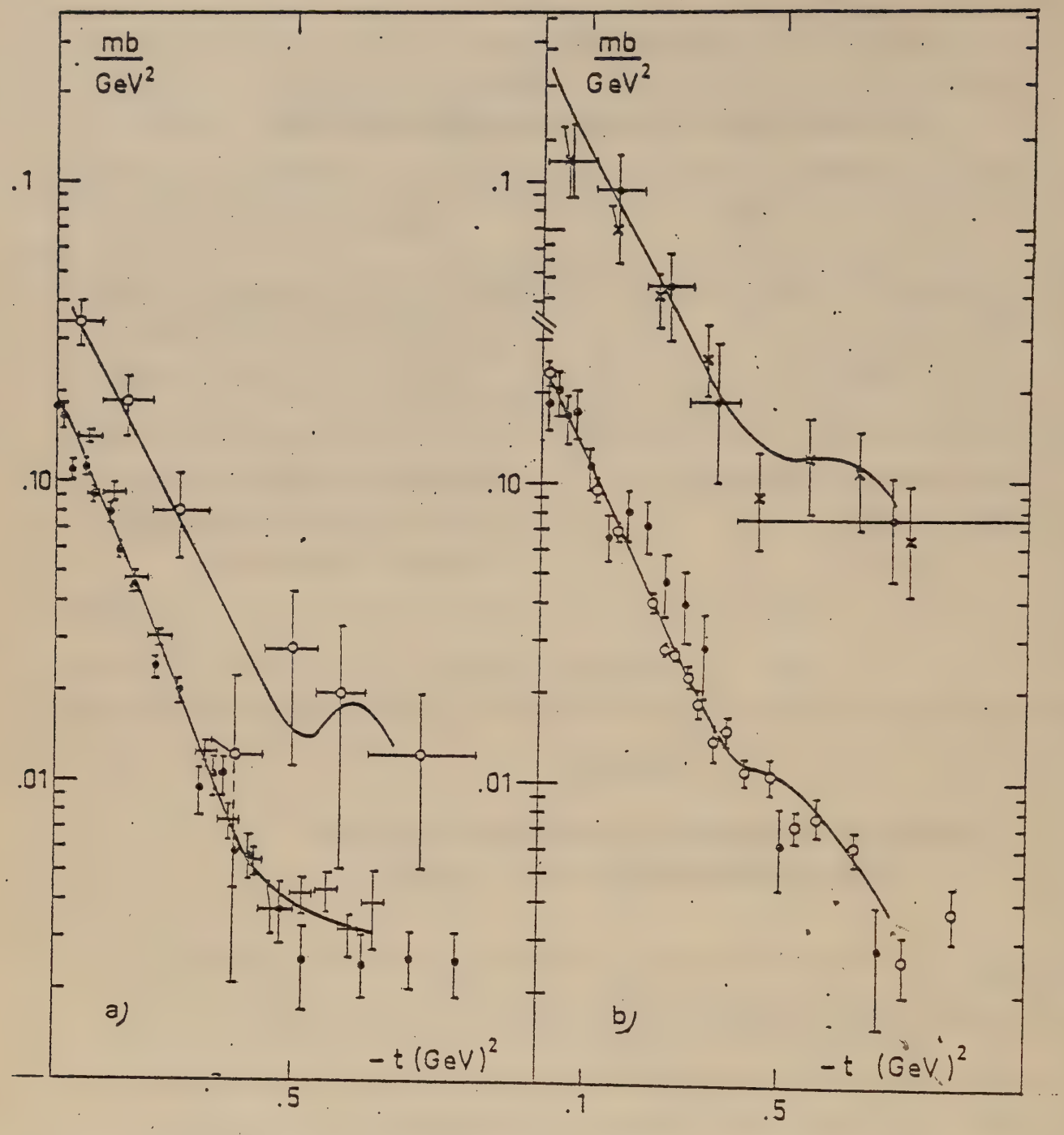
SYSTEMATICS FOR THE VECTOR AND TENSOR EXCHANGES

	Vectors : ρ, ω, K_V^*	Tensors : A_2, f_o, K_T^*.
n=0	Peripheral zero in the Im part. absorbed Regge pole or $J_o(R\sqrt{-t})$ with $R \sim 1$fm	Compatible with a Regge pole with a no compensation mechanism : -Simple zero in ReT -Double zero in ImT at $t \sim -.6$ GeV^2
n=1	Pure Regge pole behaviour with its WSNZ or $J_1(R\sqrt{-t})$	Regge Pole behaviour, with an approximate exchange degeneracy with its vector partner.
n=2	Not yet well known.	

TABLE II

VALUES OF THE DIFFERENT COUPLING FACTOR TOGETHER WITH THE CORRESPONDING SU(3) F COUPLINGS.

+	a_V	a_T	b_V	b_T
Σ^+	$1.22 \pm .05$	$1.19 \pm .1$	$-.45 \pm .25$	$-.34 \pm .05$
Λ^o	$1.16 \pm .13$	1.03 ± 12	$.45 \pm .03$	$.43 \pm .05$
F coupling (F+D =1)	1.26	1.39	.21	.26
F values from [4]	$F_+ \simeq 1.5$		$F_- \simeq 0.25$	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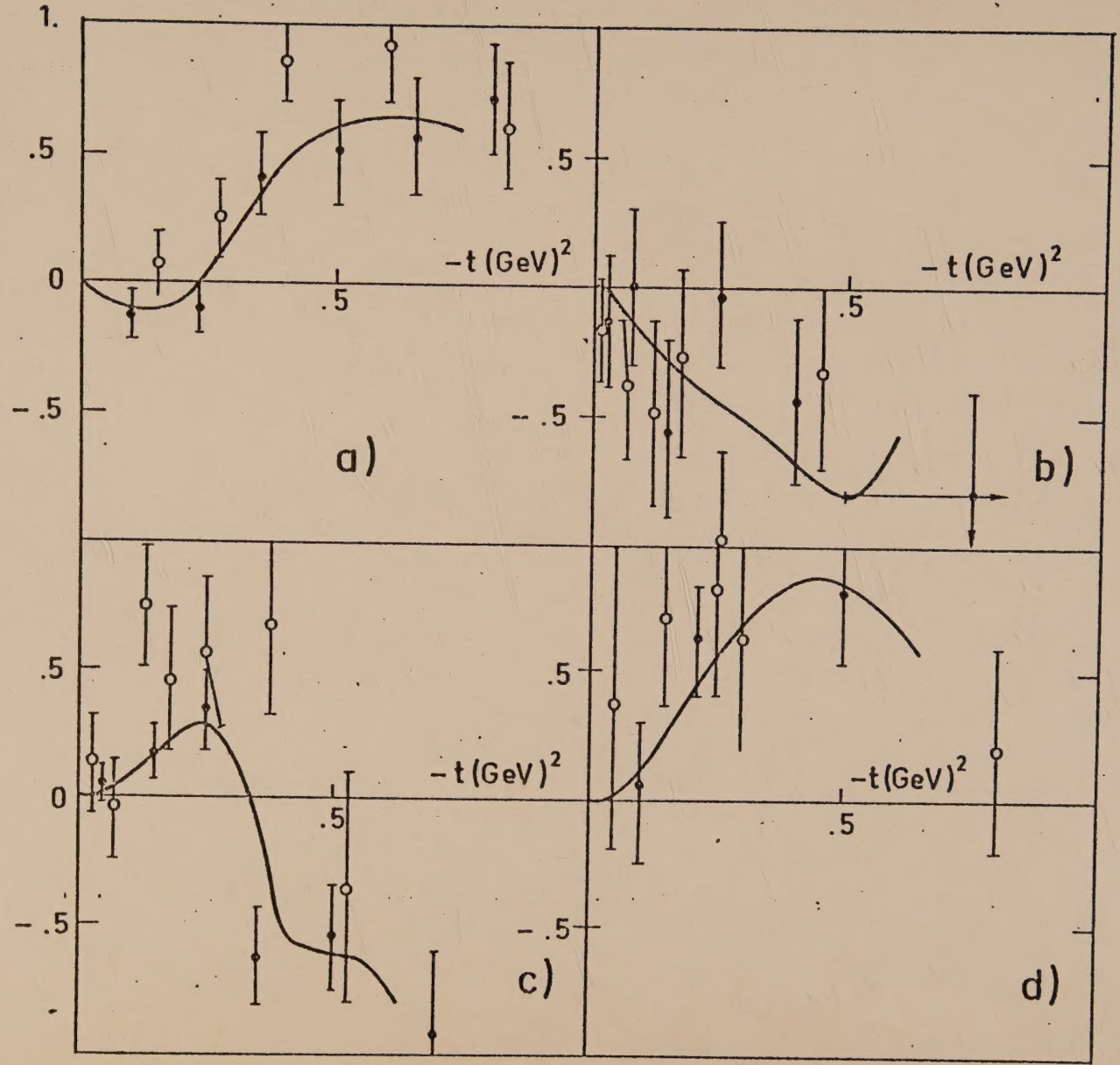
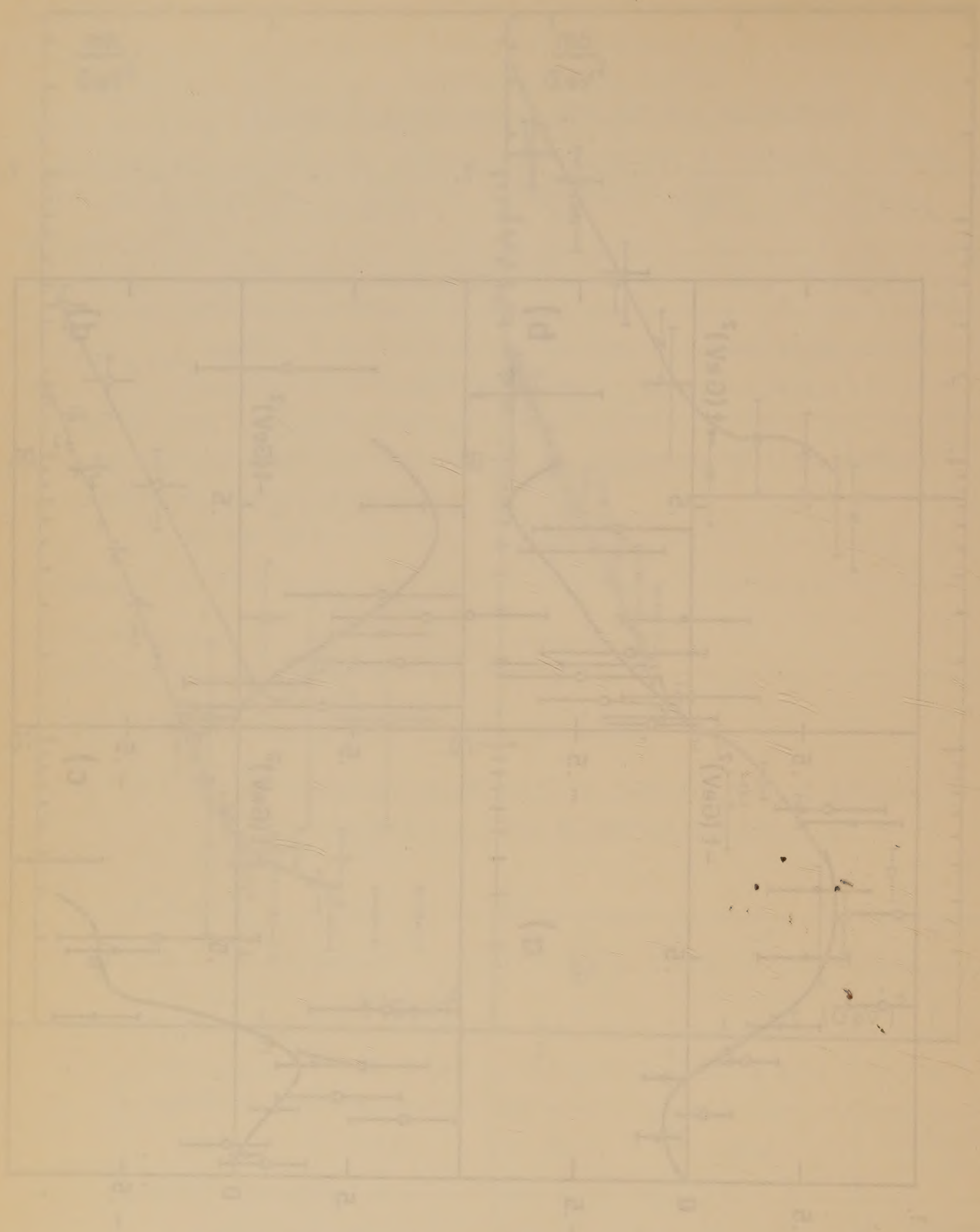
GeV²
.1
.10
.01
a)
-t (GeV)²
.5
mb
GeV²
.1
.10
.01
b)
-t (GeV)²
.1
.5


1.
.5
0
-.5
-t (GeV)²
.5
a)
b)
c)
d)